Reflection and absorption of a high-power ultrashort laser pulse heating a solid-state target*

V.A. ISAKOV, A.P. KANAVIN, AND S.A. URYUPIN

P. N. Lebedev Physical Institute, Russian Academy of Sciences, Moscow, Russia

Abstract

Patterns of absorption and reflection of an ultrashort laser pulse by plasma with a solid-state density, which are realized in the mode of the normal skin effect, are studied. If the pulse power is so high to provide the electron heating during the time shorter than the reciprocal fundamental frequency, a nonlinear suppression of absorption takes place and the spectrum of the reflected radiation contains odd harmonics of the fundamental frequency.

Keywords: Absorption coefficient; Dense plasma; Femtosecond laser pulses; Normal skin affect; Odd harmonics

1. INTRODUCTION

In the last few years, significant progress in ultrashort pulse generation (Batani & Wooton, 2004; Fuerbach *et al.*, 2005) has obtained a pulse duration τ_p compared to an oscillation period corresponding to a fundamental frequency ω of laser radiation (Zavelani-Rossi *et al.*, 2002; Schenkel *et al.*, 2003). Such short pulses react to a solid target by producing plasma at the target surface with a solid-state density, $n \approx 10^{23}$–10^{24} cm^{-3}, moreover, if the pulse power is relatively low, the electron temperature is also low. If this is the case, a hydrodynamic expansion of the plasma and the corresponding change of the density are negligible at the time scale $\sim\tau_p$, that is, one can assume an abrupt plasma interface. Since the plasma density is high, the frequency ν of electron collisions at an electron temperature of tens of electron-volts significantly exceeds the fundamental frequency of radiation ω and the mean free electron path l is much smaller than the depth of the skin layer d that corresponds to conditions of a normal skin effect. The electron heating is caused by collisions with ions and their cooling is caused by heat removal deep into the plasma, moreover, due to condition $l \ll d$, the heat transport can be described within the Spitzer–Harm approximation. Such representations of the ultrashort laser pulse interaction with a dense plasma are on the basis of several theoretical researches (Rozmus & Tikhonchuk, 1990; Rozmus *et al.*, 1996; Honrubia & Tikhonchuk, 2004; Isakov *et al.*, 2004), mainly devoted to the studies of the

influence of the electron temperature evolution in the skin layer on the laser radiation absorption and reflection, as well as to study the heat removal from the skin layer. In this connection, the conditions where the electron temperature changes during the time above the period of the light field oscillations $\sim 1/\omega$ are under consideration, and the pulse duration is so high that an influence of effects arising due to its finiteness can be eliminated. At the same time, an inverse situation is possible where the electron heating occurs during the time shorter than $\sim 1/\omega$ and the pulse duration is commensurable to $\sim 1/\omega$. Under such conditions, first it is necessary to take into account the change of the electron temperature in the course of the field penetration into the plasma, and second the moments of the laser pulse switching-on and switching-off must be taken into account. An account for the influence of these effects on the patterns of a high-power ultrashort laser pulse interaction with dense plasma is the aim of this paper.

2. GENERAL RELATIONSHIPS

Let us consider the interaction of an electromagnetic pulse with plasma occupying the space region $z > 0$. We assume that the pulse normally incidences on the plasma surface and the distribution of its electric and magnetic fields is described by relationships

$$\mathbf{B}_i(z,t) = (0, B_i(z,t), 0), \quad \mathbf{E}_i(z,t) = (E_i(z,t), 0, 0), \tag{1}$$

$$B_i(z,t) = E_i(z,t) = E_0 \exp\left[-\frac{(\omega t - kz)^2}{\omega^2 \tau_p^2}\right] \sin(\omega t - kz), \tag{2}$$

Address correspondence and reprint requests to: S.A. Uryupin, P.N. Lebedev Physical Institute, Leninskii Pr. 53, Moscow 119991, Russia. E-mail: uryupin@sci.lebedev.ru
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where E_0 is the strength of the electric field, $\mathbf{k} = (0,0,k)$ is the wave vector, $\omega = kc$ is the fundamental radiation frequency, c is the speed of light, and τ_p is the pulse duration. Fields of type (1) and (2) satisfy the Maxwell equations in vacuum. The electromagnetic pulse is partially reflected by the plasma. According to the Maxwell equations, the reflected field is of the form

$$\mathbf{E}_r(z,t) = (E_r(z + ct),0,0),$$
$$\mathbf{B}_r(z,t) = (0, B_r(z + ct),0), = (0,-E_r(z + ct),0). \quad (3)$$

The incident and reflected fields are connected with the fields $\mathbf{E}$ and $\mathbf{B}$ inside the plasma, $\mathbf{E}(z,t) = (E(z,t),0,0)$ and $\mathbf{B}(z,t) = (0, B(z,t),0)$. The electric field inside the plasma significantly depends on the form of connection between the current density $j(z,t)$ and the field $E(z,t)$. Further we consider, sufficiently dense hot plasmas where the electron plasma frequency $\omega_L = (4\pi e^2 n/m)^{1/2}$ is much higher than the frequency of electron–ion collisions $\nu = 4\pi Z e^4 n\Lambda m^{-2} v_T^{-3}$, where e and m are the electron charge and mass, n and v_T are the electron density and thermal velocity, Λ is the Coulomb logarithm, and Z is the ionization degree of ions. For a wide class of hot plasmas $Z \gg 1$ that is further assumed. We also assume that the collision frequency ν is much higher than typical frequencies of changing the electron temperature T and electric field E, and a mean free path of thermal electrons $\ell = v_T/\nu$ are much less than typical scales of the field and temperature inhomogeneities. Besides, let us assume that the electric field strength is rather small, providing the following inequality: $\sqrt{Z}E \ll m\nu v_T|e|^{-1}$. Within the framework of the assumptions made, a local in space and time linear relation between the current density and the field takes place:

$$\mathbf{j}(z,t) = \frac{32}{\sqrt{2\pi}} \frac{e^2 n}{m\nu} \mathbf{E}(z,t) \equiv \frac{4\sqrt{2}}{\pi\sqrt{\pi}} \frac{\omega_L^2}{\nu} \mathbf{E}(z,t). \quad (4)$$

Using relation (4) for the field in the plasma, one has the following approximate equation:

$$\frac{c^2}{\omega_L^2} \frac{\partial^2}{\partial z^2} E(z,t) = \frac{32}{\sqrt{2\pi}} \frac{\partial}{\partial t}\left[\frac{1}{\nu} E(z,t)\right]. \quad (5)$$

Eq. (5) should be supplemented by an initial and two boundary conditions. The initial condition corresponds to the fact that for $t \to -\infty$ there is no electromagnetic field in the plasma,

$$E(z, t \to -\infty) = 0, \quad z \geq 0. \quad (6)$$

The tangential components of the electric and magnetic fields are continuous at the plasma boundary. Taking this into account, one has with the same accuracy as for Eq. (5)

$$\left.\frac{\partial}{\partial z} E(z,t)\right|_{z=0} = \left(\frac{4t}{\tau_p^2} \sin\omega t - 2\omega\cos\omega t\right) \frac{E_0}{c} \exp\left(-\frac{t^2}{\tau_p^2}\right). \quad (7)$$

The second boundary condition corresponds to the absence of the field deep in the plasma,

$$E(z \to \infty, t) = 0. \quad (8)$$

The frequency of electron–ion collisions, entering Eq. (5), depends on the temperature. Since $\partial \ln \nu/\partial t = -(3/2)\partial \ln T/\partial t$, the changing in time of the collision frequency is determined by the corresponding changing in time of the electron temperature. The evolution of the electron temperature in the plasma is described by the following equation

$$\frac{3}{2} nk_B \frac{\partial}{\partial t} T(z,t) = Q(z,t) - \frac{\partial}{\partial z} q(z,t), \quad (9)$$

where k_B is the Boltzmann constant, $Q(z,t)$ is the Joule heat released at point z and at time t,

$$Q(z,t) = j(z,t)E(z,t) = \frac{4\sqrt{2}}{\pi\sqrt{\pi}} \frac{\omega_L^2}{\nu} E^2(z,t), \quad (10)$$

$q(z,t)$ is the heat-flux density deep in the plasma,

$$q(z,t) = -\lambda(z,t) \frac{\partial}{\partial z} T(z,t), \quad (11)$$

$$\lambda(z,t) = \frac{128}{\sqrt{2\pi}} nk_B \frac{v_T^2}{\nu} \equiv \frac{128}{\sqrt{2\pi}} nk_B l v_T, \quad (12)$$

is the Spitzer–Harm factor of heat conductivity. Note that in conditions under consideration an increase in temperature is insignificant. That is why the ratio of the mean free electron path to the typical scale of temperature inhomogeneity is relatively low during the period of laser pulse action. Thus above-thermal electrons cannot influence the energy removal from the skin layer to an extent providing significant deviations from the Spitzer–Harm law and values of temperature in the skin layer.

There is the following initial condition for Eq. (9):

$$T(z, t \to -\infty) = T_0, \quad (13)$$

where T_0 is the electron temperature before the laser pulse action, as well as two boundary conditions,

$$T(z \to \infty, t) = T_0, \quad \left.\frac{\partial}{\partial z} T(z,t)\right|_{z=0} = 0. \quad (14)$$

The boundary condition at $z = 0$ means the absence of the heat flux at the plasma surface. Eqs. (5) and (9) allows one to study the evolution of the field distribution and electron temperature in the dense hot plasma under the action of an ultrashort laser pulse.

3. HEATING OF ELECTRONS AND NONLINEAR FIELD PENETRATION INTO THE PLASMA

In the case of the action of laser pulses of high power inequality $I = cE_0^2/8\pi \gg nk_BTc/4$ can be fulfilled. If this inequality is satisfied and heat losses from the skin layer are low, the time $1/\omega$ is not short as compared to a typical time of the electron temperature change. Under such conditions, it is necessary to construct a solution for Eq. (5) taking into account a rather fast evolution of the electron temperature described by Eq. (9). Obtaining a numerical solution for Eqs. (5) and (9), let us introduce dimensionless time- and space-variables $\tau = \omega t$, $\xi = z/d_0$, and dimensionless functions $y = T/T_0$ and $u = E/E(0)$ that describe the temperature and field distribution inside the plasma. Here, the following notation is used:

$$d_0 = \frac{(2\pi)^{1/4}}{4} \frac{c}{\omega_L} \sqrt{\frac{\nu_0}{\omega}}, \quad E(0) = 2\frac{\omega}{c} d_0 E_0, \quad (15)$$

where $\nu_0 = \nu(T_0)$ is the effective frequency of electron–ion collisions at temperature T_0. Using these variables, one obtains from (5) and (9)

$$\frac{\partial}{\partial \tau}[u(\xi,\tau)y^{3/2}(\xi,\tau)] = \frac{1}{2}\frac{\partial^2}{\partial \xi^2} u(\xi,\tau), \quad (16)$$

$$\frac{\partial}{\partial \tau} y(\xi,\tau) = \frac{16}{3} Wu^2(\xi,\tau)y^{3/2}(\xi,\tau)$$
$$+ \frac{16}{3} D \frac{\partial}{\partial \xi}\left[y^{5/2}(\xi,\tau)\frac{\partial}{\partial \xi}y(\xi,\tau)\right], \quad (17)$$

where the following notation is used:

$$W = \frac{E_0^2}{4\pi nk_BT_0} \equiv 2\frac{I}{cnk_BT_0}, \quad (18)$$

$$D = \frac{16}{\sqrt{2\pi}}\frac{\nu_0}{\omega}\left(\frac{l_0}{d_0}\right)^2 \equiv \left(\frac{16}{\sqrt{2\pi}}l_0\frac{\omega_L}{c}\right)^2, \quad (19)$$

and $l_0 = \sqrt{k_BT_0/m}/\nu_0$ is the mean free path at temperature T_0. In this case, initial (6) and (13) and boundary (7), (8), and (14) conditions for functions $u(\xi,\tau)$ and $y(\xi,\tau)$ are of the form

$$u(\xi \geq 0, \tau \to -\infty) = 0, \quad y(\xi \geq 0, \tau \to -\infty) = 1, \quad (20)$$

$$\left.\frac{\partial}{\partial \xi}u(\xi,\tau)\right|_{\xi=0} = \left(\frac{2\tau}{T_p^2}\sin\tau - \cos\tau\right)\exp\left(-\frac{\tau^2}{T_p^2}\right),$$
$$u(\xi \to \infty, \tau) = 0, \quad (21)$$

$$\left.\frac{\partial}{\partial \xi}y(\xi,\tau)\right|_{\xi=0} = 0, \quad y(\xi \to \infty, \tau) = 1, \quad (22)$$

where $T_p = \omega\tau_p$. Numerical solution for Eqs. (16) and (17) for $T_p = 2\pi$, $D = 10^{-3}$, and $W = 0.1, 0.5, 2$ are presented in

Figures 1 and 2. Figure 1 shows the space distribution of the squared field in the plasma at the time moment $\tau = 0$ for three values of $W = 0.1, 0.5$, and 2. The electron temperature time-evolution at the plasma surface ($\xi = 0$) for $W = 0.1$, 0.5, and 2 is shown in Figure 2. As is seen from the dependences, the electron temperature increases and the squared field at the plasma surface decreases with an increase in the radiation flux density, that is., with an increase in the parameter W. The space region of the field concentration decreases with an increase in W. The decreases in the effective depth of the skin layer and the modulus of the field at the surface are due to fast electron heating.

Cooling of electrons due to the heat removal deep into the plasma is characterized by the parameter D and can be seen in Figure 2. As a result of the electron heating, the effective value of D increases providing the appearance of dips of the curves in Figure 2 for time evolution of the electron temperature at the plasma surface. The greater W, the stronger the electron heating and the deeper the dips of the curves $y(\xi = 0, \tau)$. In the course of time, at the laser pulse shutdown, the electron heating becomes slower but the energy removal from the skin layer is still high that results in a monotonic decrease in the electron temperature at the surface (see Fig. 2).

4. ABSORPTION OF A HIGH-POWER PULSE AND HARMONIC GENERATION

Changes in the features of the high-power laser pulse penetration into the plasma caused by the fast heating of electrons result in changes in the absorption coefficient and the spectrum of the reflected radiation. Taking into account the explicit form of the incident pulse field (2), one can obtain the following expression for the absorption coefficient:
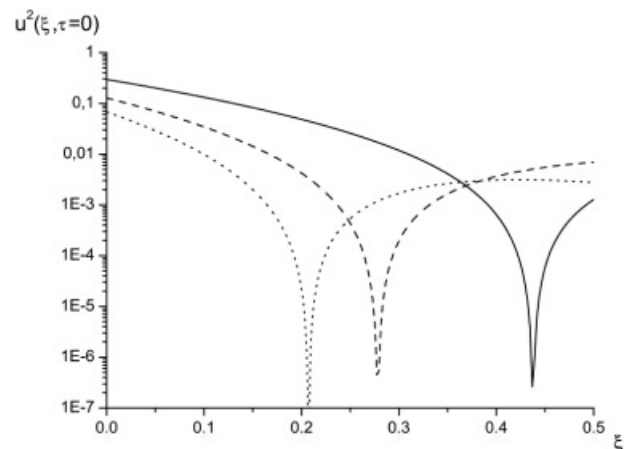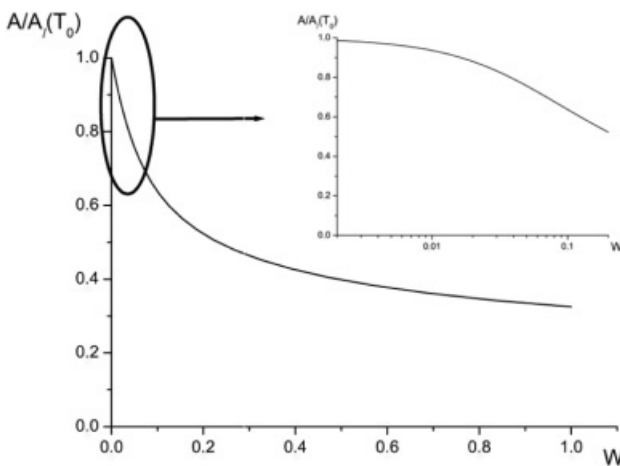

Fig. 1. Distribution of the squared field in the plasma at the time moment $\tau = 0$. The curves are obtained for $D = 10^{-3}$ and $T_p = 2\pi$ and correspond to three values of the parameter $W = 0.1$ (solid curve), 0.5 (dashed curve), and 2 (dotted curve).
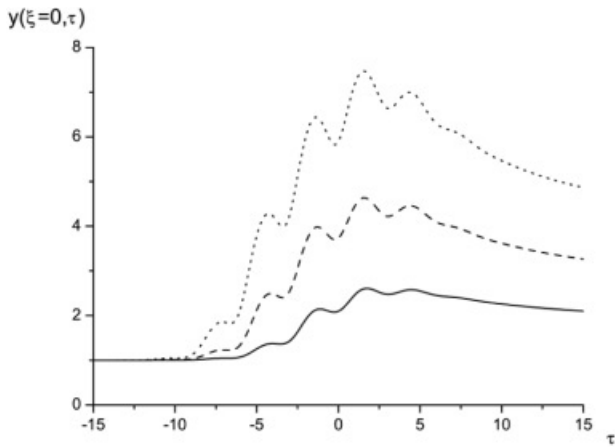
Fig. 2. Time dependence of the electron temperature at the plasma surface. Calculations are carried out for $T_p = 2\pi$, $D = 10^{-3}$, and $W = 0.1$ (solid curve), 0.5 (dashed curve), and 2 (dotted curve).

$$A = A_l(T_0)4\sqrt{\frac{2}{\pi}}\left[1 - \exp\left(-\frac{T_p^2}{2}\right)\right]^{-1}$$

$$\times \int_{-\infty}^{\infty} \frac{d\tau}{T_p} u(\xi = 0, \tau)\sin\tau \exp\left(-\frac{\tau^2}{T_p^2}\right), \quad (23)$$

where $A_l(T_0) = (2\pi)^{1/4}\sqrt{\nu_0\omega}/2\omega_L$ is the absorption coefficient for a long ($T_p \gg 1$) low-power pulse at temperature T_0. Calculation results for the absorption coefficient for a pulse with $T_p = 2\pi$ and $D = 10^{-3}$ are presented in Figure 3. According to the inset, in the case of a low radiation flux density where the heating of electrons is low, the absorption coefficient A (23) is close to $A_l(T_0)$. As the flux density increases the electron heating takes place during the action of a pulse with $T_p = 2\pi$. The fast heating is accompanied by the field reduction in the plasma and a decrease in the effective depth of field penetration. This results in a decrease

in the dissipated fraction of the laser pulse energy, and therefore, in the absorption coefficient. Figure 3 shows a nonlinear reduction in the absorption coefficient caused by the increasing parameter W. At large W the suppression of the absorption coefficient becomes slower. The curve $A(W)$ in Figure 3 is of a variable bend. This is connected with the fact that at large W, the electron temperature significantly arises to the initial stage of the laser pulse action. Such an increase in the temperature leads to a strong increase in the heat flux deep into the plasma and to a slower growth of the temperature itself. The major part of the laser pulse is absorbed under the conditions of slow changing of the electron temperature that results in a relatively slower decrease in A with the increasing W.
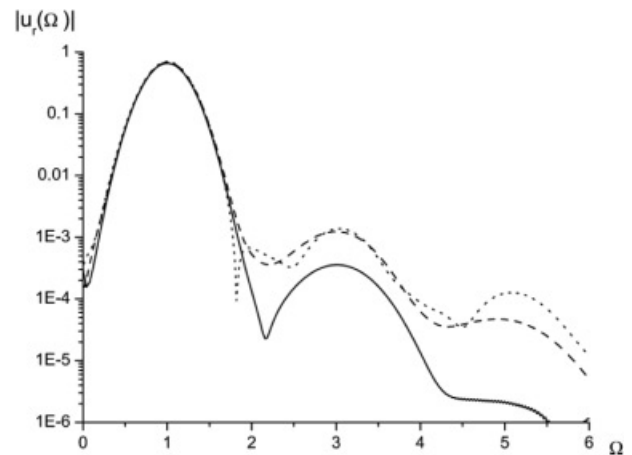
A field of the reflected pulse is identically connected to the field in the plasma. Taking into account notation (15), one obtains from the continuity condition for the electric tangential component

$$u_r(\xi = 0, \tau) = u(\xi = 0, \tau) - A_l^{-1}(T_0)\exp\left(-\frac{\tau^2}{T_p^2}\right)\sin\tau, \quad (24)$$

where $u_r(\xi = 0, \tau) = E_r(c\tau/\omega)/E(0)$. The spectral composition of the reflected radiation is determined by the Fourier component of the function $u_r(\xi = 0, \tau)$,

$$u_r(\Omega) = \int_{-\infty}^{\infty} d\tau \exp(-i\Omega\tau)u_r(\xi = 0, \tau), \quad (25)$$

where Ω is dimensionless related to the frequency ω. Figure 4 shows plots of the function $|u_r(\Omega)|$ obtained for $T_p = 2\pi$, $D = 10^{-3}$, $A_l(T_0) = 0.1$, and three values of $W = 0.01$, 0.1, and 1. As the absorption coefficient $A_l(T_0)$ is small, the reflected radiation is mainly concentrated near the frequency ω. The width of the reflection band near the fre-

Fig. 3. Absorption coefficient of a pulse with $T_p = 2\pi$ as a function of the radiation flux density W. Calculation is carried out for $D = 10^{-3}$.

Fig. 4. Modulus of the Fourier component of the field of the reflected pulse as a function of the dimensionless (related to ω) frequency Ω. Calculations are carried out for $T_p = 2\pi$, $D = 10^{-3}$, $A_l(T_0) = 0.1$, and three values of $W = 0.01$ (solid curve), 0.1 (dashed curve), and 1 (dotted curve).

quency ω is determined by the reciprocal pulse duration. Broadening of the main band due to the electron heating is relatively small, that is why the curves in Figure 4 corresponding to a different W practically coincide each other in the vicinity of $\Omega \sim 1$. The time change of the temperature causes a weak reflection at frequencies above ω. According to Eq. (17), the energy absorbed by electrons is proportional to the squared strength of the field that is concentrated mainly near the frequency ω. As a result, the temperature is modulated in time with a characteristic frequency 2ω. The electron temperature modulation causes the generation of odd harmonics of the field at frequencies $(2n + 1)\omega$, $n = 1, 2 \ldots$ Amplitude splashes of the Fourier component of the field near the frequencies 3ω and 5ω are seen in Figure 4.

Generation of harmonics of a high-power high-frequency field due to oscillations of the effective electron–ion collision frequency has been known for a long time (Silin, 1965). In Ferrante *et al*. (2003), the theory of this phenomenon was developed as applied to conditions of the normal skin effect, when $\omega_L > \nu > \omega$, the field is low-power, and the electron heating occurs during the time much longer than $1/\omega$. In Ferrante *et al*. (2003), the influence of electron-electron collisions on the oscillating motion of electrons is assumed weak, that is, valid for $\omega > \nu_{ee} = \nu/Z$, where ν_{ee} is the frequency of electron–electron collisions. If $\nu_{ee} > \omega$, the frequent electron–electron collisions maintain the electron distribution close to the Maxwellian one. Under such conditions, for description of harmonic generation, one can not use a kinetic equation but a simpler equation for the electron temperature. Within a such approach, as is shown above, the odd harmonics arise due to the fact that the frequency of electron–ion collisions is dependent on the thermal velocity modulated at frequencies $2n\omega$, $n = 1, 2, \ldots$ High-order harmonic generation by intense femtosecond laser pulses irradiated at dense plasma surfaces is also discussed by Foldes *et al*. (2003).

5. CONCLUSIONS

From the presentation above, it follows that there are specific features in optical properties of dense plasma interacting with an ultrashort laser pulse caused by fast heating of electrons. Let us illustrate the conditions where the anomalies in optical properties established above can manifest themselves. We assume as an estimate that the plasma consists of fully ionized beryllium atoms with $Z = 4$ and electrons with a density $n = 5 \cdot 10^{23}$ cm^{-3} and temperature $T_0 = 60$ eV and the fundamental frequency in the pulse is $\omega = 2 \cdot 10^{15}$ s^{-3}. With such parameters of the plasma and radiation one obtains $\Lambda \sim 1$, $\nu_0 \sim \omega_L \cong 4 \cdot 10^{16}$ s$^{-1} \gg \omega$, $l_0 \cong 10^{-8}$ cm $\ll d_0 \cong 1.3 \cdot 10^{-6}$ cm, and $2\pi/\omega \cong 3$ fs. The parameter W (18) determining a degree of the influence of fast heating of electrons on optical properties of the plasma changes from 0.04 to 1 with an increase in the radiation flux density I from $2 \cdot 10^{15}$ to $5 \cdot 10^{16}$ W/cm^2. As is seen in Figure 3, an increase in W from 0.04 to 1 result in a relative reduction in the absorption coefficient from 0.8 to 0.3.

REFERENCES

BATANI, D. & WOOTON, A.J. (2004). The international conference on ultrashort high-energy radiation and matter. *Laser Part. Beams* **22**, 197.

FERRANTE, G., ZARCONE, M. & URYUPIN, S.A. (2003). Plasma emission at the laser tripled frequency. *Phys. Lett. A* **315**, 378–383.

FOLDES, I.B., KOCSIS, G., RACZ, E., SZATMARI, S. & VEREZ, G. (2003). Generation of high harmonics in laser plasmas. *Laser Part. Beams* **21**, 517–521.

FUERBACH, A., FERNANDEZ, G., APOLONSKI, A., FUJI, T. & KRAUSZ, F. (2005). Chirped-pulse oscillators for the generation of high-energy femtosecond laser pulses. *Laser Part. Beams* **23**, 113–116.

HONRUBIA, J. & TIKHONCHUK, V.T. (2004). Workshop on simulation of ultra intense laser beams interaction with matter. *Laser Part. Beams* **22**, 95.

ISAKOV, V.A., KANAVIN, A.P. & URYUPIN, S.A. (2004). Absorption of ultrashort laser pulses heating a dense plasma. *J. Russ. Laser Res.* **25**, 156–168.

ROZMUS, W. & TIKHONCHUK, V.T. (1990). Skin effect and interaction of short laser pulses with dense plasma. *Phys. Rev. A* **42**, 7401–7412.

ROZMUS, W., TIKHONCHUK, V.T. & CAUBLE, R. (1996). A model of ultrashort laser pulse absorption in solid targets. *Phys. Plasmas* **3**, 360–371.

SCHENKEL, B., BIEGERT, J., KELLER, U., VOZZI, C., NOSOLI, M., SANSONE, G., STAGIRA, S., DE SILVESTRI, S. & SVELTO, O. (2003). Generation of 3.8-fs pulses from adaptive compression of a cascaded hollow fiber super continuum. *Opt. Lett*. **28**, 1987–1996.

SILIN, V.P. (1965). Nonlinear high-frequency plasma conductivity. *JETP* **20**, 1510 .

ZAVELANI-ROSSI, M., POLLI, D., CERULLO, G., DE SILVESTRI, S., GALLMANN, L., STEINMEYER, G. & KELLER, U. (2002). Few-optical-cycle laser pulses by OPA: broadband chirped mirror compression and SPIDER characterization. *Appl. Phys. B* **74**, S245.