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Collisional absorption of dense plasmas in strong laser fields

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Abstract

Collisional absorption of dense fully ionized plasmas in strong laser fields is investigated using quantum kinetic theory as well as molecular dynamics simulations. Quantum statistical calculations are presented for the important case of a two-temperature plasma. Comparision is given to molecular dynamics simulations. Special attention is devoted to the question of how to treat properly the attractive electron–ion interaction for short ranges in classical molecular dynamics simulations.

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1. Introduction

An important question in almost all experiments with interaction of intense laser pulses with matter is the calculation of the energy deposition and the description of the heating connected with that. If a solid target is irradiated by such an intense laser pulse, dense plasmas can be created relevant for astrophysical research and for inertial confinement fusion.

Especially, at high intensities the quiver velocity can be large compared to the thermal velocity and interesting nonlinear effects have to be expected. One of the important mechanisms of energy deposition is collisional absorption usually described in terms of the electron–ion collision frequency [1–4].

2. Collisional absorption in fully ionized plasmas

In the case of strong fields, the dependence of, e.g., the electrical current density on the electric field has an exponential form and causes thus nonlinear effects such as multi-photon absorption

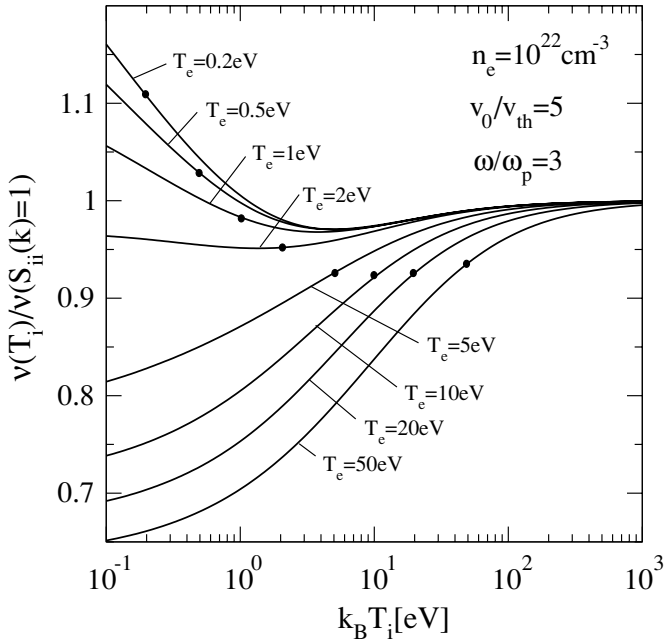

Figure 1. Ratio of collision frequencies with HNC structure factor to the respective values with $S_{ii} = 1$ as a function of the ion temperature T_i. There are shown curves for different electron temperatures T_e; as a guide for the eye, the circles mark the case of an isothermal plasma where $T_e = T_i$.

and the occurrence of higher harmonics. For a harmonic electric field, $\mathbf{E} = \mathbf{E}_0 \cos \omega t$, the cycle averaged dissipation of energy is given by [5–7]

$$\langle \mathbf{j} \cdot \mathbf{E} \rangle = n_i \int \frac{d^3q}{(2\pi\hbar)^3} \frac{V_{ei}^2(q)}{V_{ee}(q)} \mathcal{S}_{ii}(\mathbf{q}, T_i) \sum_{n=-\infty}^{\infty} n\omega \, J_n^2 \left(\frac{\mathbf{q} \cdot \mathbf{v}_0}{\hbar\omega} \right) \mathrm{Im}\, \varepsilon_{ee}^{-1}(\mathbf{q}, -n\omega, T_e), \tag{1}$$

with ε_{ee} being the dielectric function of the electron component and $\mathcal{S}_{ii}$ the static structure factor of the ions. The sum over Bessel functions J_n reflects the nonlinear dependence on the field; $v_0 = eE_0/(m_e\omega)$ is the quiver velocity of free electrons. With equation (1), collisional absorption can be described for the case of two-temperature plasmas. Also molecular dynamics (MD) simulations confirm [6] that mainly the electrons are heated by the laser field whereas the temperature of the ions stays almost constant. The influence of the ion component with temperature T_i is accounted for in equation (1) by the static ionic structure factor. Results are shown in figure 1 for an ion structure factor in hypernetted chain (HNC) approximation. The calculations, given here for $v_0/v_{th} = 5$ (for smaller fields see [6]), show a considerable influence of structure factor effects on collisional absorption especially for the case $T_i < T_e$. Only for comparatively low electron temperatures, there is an enhancement of the collision frequency at all.

3. Molecular dynamics simulations

The energy absorption was also calculated with molecular dynamics simulations [6, 8]. The external electric field was implemented as a homogeneous linearly polarized harmonic field. The main difficulty in order to simulate a fully ionized plasma is to model the attractive

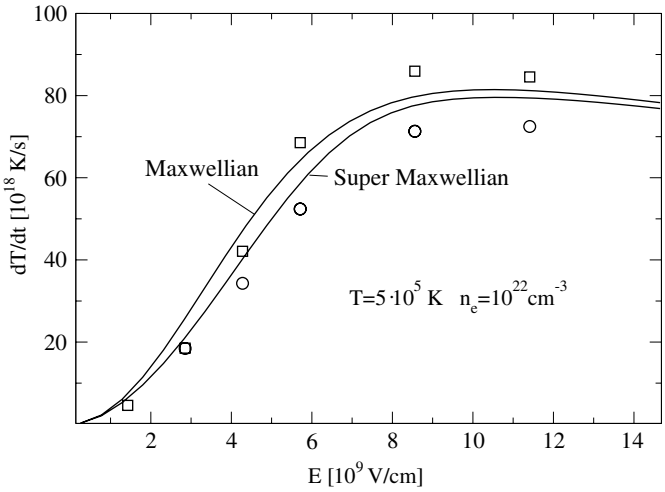

Figure 2. Heating rate of the electrons in a two-temperature hydrogen plasma ($T_i = 1000$ K) as a function of the applied field strength. The solid lines denote the analytic results for two different assumptions for the velocity distribution functions. The circles are MD results for the Deutsch potential with a thermal wavelength, the squares are results for the Deutsch potential with De Broglie wavelength with inclusion of quiver motion.

electron–ion interaction: the pure Coulomb potential has a singularity at the origin which causes a non-physical behaviour of the system. The simplest way to avoid this divergence is to use a so-called soft Coulomb potential [9] by introduction of some arbitrary cut-off at small distances. Another possibility is the Kelbg potential [10]

$$\Phi_{ij} = \frac{q_i q_j}{4\pi\epsilon_0 r}\left\{1 - \exp\left(-r^2/\lambda_{ij}^2\right) + \sqrt{\pi}r/\lambda_{ij}[1 - \mathrm{erf}(r/\lambda_{ij})]\right\}, \tag{2}$$

which was derived from equilibrium quantum statistics using the two-particle Slater sum. This potential is temperature dependent via the thermal wavelength $\lambda_{ij} = \hbar/\sqrt{2\mu_{ij}k_B T_{ij}}$, where μ_{ij} denotes the reduced mass. For the temperature in λ_{ei}, we used $T_{\mathrm{ei}} \approx T_{\mathrm{e}}$. A quite similar potential is the frequently used Deutsch potential [11]. In strong electric fields, the assumption of local thermodynamic equilibrium may not be fulfilled, there are attempts to use the DeBroglie wavelength including the quiver velocity instead of the thermal wavelength [8, 12]. The validity of this approach still needs some further discussion.

The temperature can be defined in MD simulations taking into account the undirected motion only. For the linearly polarized laser field under consideration, there occurs however an anisotropy. One can define longitudinal (in the field direction) and transversal temperatures according to

$$\frac{1}{2}k_B T_{\mathrm{a}\parallel} = m_{\mathrm{a}}\frac{\langle v_{\mathrm{a}\parallel}^2\rangle - \langle \mathbf{v}_{\mathrm{a}\parallel}\rangle^2}{2}, \qquad k_B T_{\mathrm{a}\perp} = m_{\mathrm{a}}\frac{\langle v_{\mathrm{a}\perp}^2\rangle}{2}, \tag{3}$$

where the angles denote an averaging over all particles of species 'a'. In addition a mean temperature of species 'a' may be defined as

$$\tfrac{3}{2}k_B T_{\mathrm{a}} = E_{\mathrm{therm}} = \tfrac{1}{2}k_B T_{\mathrm{a}\parallel} + k_B T_{\mathrm{a}\perp}. \tag{4}$$

The MD calculations were performed using periodic boundary conditions with Ewald summation. The number of particles was between 2000 and 5000. In contrast to Pfalzner and

Gibbon who used the so-called constant temperature dynamics [9] in order to determine the heating rate, we preferred a dynamical method because it simulates the real heating process instead of using an artificial heat bath. The change of thermal energy is associated with the electron–ion collision frequency for the high-frequency case via

$$\nu_{\mathrm{ei}}(\omega) = \frac{\omega^2}{\omega_{\mathrm{p}}^2}\frac{2}{\epsilon_0 E_0^2}\langle \mathbf{j}\cdot\mathbf{E}\rangle = \frac{\omega^2}{\omega_{\mathrm{p}}^2}\frac{2n_{\mathrm{e}}}{\epsilon_0 E_0^2}\frac{\mathrm{d}E_{\mathrm{therm}}}{\mathrm{d}t}. \tag{5}$$

For details of the calculations, see [6, 8]. Longitudinal and transversal temperatures differ from each other, details will be presented elsewhere. Usually, the longitudinal temperature is larger than the transversal one. For high field strengths, however, this relation turns into the opposite: the width of the velocity distribution function in field direction becomes smaller than for a Maxwellian [6].

Figure 2 shows the heating of electrons at a mean temperature $T_{\mathrm{e}} = 5 \times 10^5$ K as a function of the applied field strength. The circles and squares denote the simulation data of the corresponding two-temperature plasma where the ion temperature was 1000 K. The solid lines denote results from the analytic approach for a two-temperature plasma with different assumptions for the velocity-distribution function: we considered Maxwellian as well as super-Maxwellian distribution functions [13]. The simulation data for the two different potentials show large deviations from each other for field strengths above 4×10^9 V cm^{-1}. A smaller effective cut-off radius enhances the probability of large-angle scattering and leads therefore to stronger electron heating. There is also an influence of the cut-off radius on the distribution function. Calculations with a soft Coulomb potential with different cut-off radii at fixed field strength show, however, no systematic trend in the deviation between longitudinal and transversal temperatures.

To summarize, the choice of the effective electron–ion interaction potential remains a crucial issue and needs further discussion.

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