Stopping of ions in a plasma irradiated by an intense laser field

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Abstract

The inelastic interaction between heavy ions and an electron plasma in the presence of an intense radiation field (RF) is investigated. The stopping power of the test ion averaged with a period of the RF has been calculated assuming that $\omega_0 > \omega_p$, where ω_0 is the frequency of the RF and ω_p is the plasma frequency. In order to highlight the effect of the radiation field we present a comparison of our analytical and numerical results obtained for nonzero RF with those for vanishing RF. It has been shown that the RF may strongly reduce the mean energy loss for slow ions while increasing it at high–velocities. Moreover, it has been shown, that acceleration of the projectile ion due to the RF is expected at high–velocities and in the high–intensity limit of the RF, when the quiver velocity of the plasma electrons exceeds the ion velocity.

Keywords: Stopping power; Inelastic interaction; Radiation field

1 Introduction

The interaction of charged particles with a plasma in the presence of radiation field (RF) has been a subject of great activity, starting with the work of Tavdgiridze, Aliev, Gorbunov and other authors (Tavdgiridze & Tsintsadze, 1970; Aliev *et al.*, 1971; Arista *et al.*, 1989; Akopyan *et al.*, 1997; Nersisyan & Akopyan, 1999). A comprehensive treatment of the quantities related to inelastic particle–solid and particle–plasma interactions, like scattering rates and differential and total mean free paths and energy losses, can be formulated in terms of the dielectric response function obtained from the electron gas model. The results have important applications in radiation and solid–state physics (Ritchie *et al.*, 1975; Tung & Ritchie, 1977; Echenique, 1987), and more recently, in studies of energy deposition by ion beams in plasma fusion targets (Arista & Brandt, 1981; Mehlhorn, 1981; Maynard & Deutsch, 1982; Arista & Piriz, 1987; Avanzo *et al.*, 1993; Couillaud *et al.*, 1994). On the other hand, the achievement of high–intensity laser beams with frequencies ranging between the infrared and vacuum–ultraviolet region has given rise to the possibility of new studies of interaction processes, such as electron–atom scattering in laser fields (Kroll & Watson, 1973; Weingartshofer *et al.*, 1977, 1983), multiphoton ionization (Lompre *et al.*, 1976; Baldwin & Boreham, 1981), inverse bremsstrahlung and plasma heating (Seely & Harris, 1973; Kim & Pac, 1979; Lima *et al.*, 1979), screening breakdown (Miranda *et al.*, 2005), and other processes of interest for applications in optics, solid–state, and fusion research.

In this paper we present a study of the effects of intense RF on the interaction of nonrelativistic particles with an electron plasma. The problem is formulated using the random–phase approximation (RPA), and includes the effects of the RF in a self–consistent way. The electromagnetic field is treated in the long–wavelength limit, and the electrons are considered nonrelativistic. These are good approximations provided that (i) the wavelength of the RF ($\lambda_0 = 2\pi c/\omega_0$) is much larger than the typical screening length ($\lambda_s = v_s/\omega_p$ with v_s the mean velocity of the electrons and ω_p the plasma frequency), and (ii) the "quiver velocity" of the electrons in the RF ($v_E = eE_0/m\omega_0$) is much smaller than the speed of light c. These conditions can be alternatively written as (i) $\omega_0/\omega_p \ll 2\pi c/v_s$, (ii) $W_L \ll \frac{1}{2}n_0 c(mc^2)(\omega_0/\omega_p)^2$, where $W_L = cE_0^2/8\pi$ is the RF intensity. As an estimate in the case of dense gaseous plasma, with electron density $n_0 = 10^{18}$ cm^{-3}, we get $\frac{1}{2}n_0 mc^3 \simeq 1.2 \times 10^{15}$ W/cm^2. Thus the limits (i) and (ii) are well above the values obtained with currently available high–power RF sources, and so the approximations are well justified.

We have calculated the effects of the RF on the mean energy loss (stopping power) of the test ion considering two somewhat distinct cases with slow and fast projectiles moving in a classical and fully degenerated electron gas,

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respectively. In the latter case the degenerated electron gas is treated within a simple plasmon–pole approximation proposed by Basbas and Ritchie (Basbas & Ritchie, 1982). It has been shown, that besides usual stopping in a plasma it is possible to accelerate the charged particles beam through RF. This effect is expected for fast projectiles and in the high–intensity limit of the RF, when the "quiver velocity" of the plasma electrons exceeds the projectile ion velocity.

2 RPA formulation

We consider the time–dependent Hamiltonian for the plasma electrons in the presence of both a radiation field (RF) with vector potential $\mathbf{A}(t) = (c/\omega_0)\mathbf{E}_0\cos(\omega_0 t)$, and a self–consistent scalar potential $\varphi(\mathbf{r}, t)$ (Arista *et al.*, 1989; Nersisyan & Akopyan, 1999), i.e.,

$$H(t) = \sum_{\mathbf{p}} \frac{1}{2m}\left(\mathbf{p} - \frac{e}{c}\mathbf{A}(t)\right)^2 c^+_{\mathbf{p}} c_{\mathbf{p}} - e\sum_{\mathbf{p},\mathbf{k}} \varphi(\mathbf{k}, t) c^+_{\mathbf{p}+\mathbf{k}} c_{\mathbf{p}}, \tag{1}$$

where $c_{\mathbf{p}}$, $c^+_{\mathbf{p}}$ are annihilation and creation operators for electrons with momentum $\mathbf{p}$, respectively, and $\varphi(\mathbf{k}, t)$ is the Fourier transform of $\varphi(\mathbf{r}, t)$.

The potential $\varphi(\mathbf{k}, t)$ is produced by the external charge and by the induced electronic density, viz.,

$$k^2\varphi(\mathbf{k}, t) = 4\pi\rho_0(\mathbf{k}, t) - 4\pi e\sum_{\mathbf{p}} N_{\mathbf{p}}(\mathbf{k}, t) \tag{2}$$

being $\rho_0(\mathbf{k}, t)$ the Fourier transform of the external charge density $\rho_0(\mathbf{r}, t)$, and $N_{\mathbf{p}}(\mathbf{k}, t) = (c^+_{\mathbf{p}-\mathbf{k}} c_{\mathbf{p}})_t$ is the electrons number operator.

The time evolution of the operator $N_{\mathbf{p}}(\mathbf{k}, t)$ is determined by the equation

$$i\hbar\frac{\partial N_{\mathbf{p}}(\mathbf{k}, t)}{\partial t} = \left[N_{\mathbf{p}}(\mathbf{k}, t), H(t)\right]. \tag{3}$$

In particular, for an oscillatory field $\mathbf{A}(t)$ and within random–phase approximation (RPA) Eq. (3) has the solution (Arista *et al.*, 1989; Nersisyan & Akopyan, 1999)

$$N_{\mathbf{p}}(\mathbf{k}, t) = \frac{ie}{\hbar}\left(f_{\mathbf{p}-\mathbf{k}} - f_{\mathbf{p}}\right)\int_{-\infty}^{t} dt'\varphi(\mathbf{k}, t')\exp\left[\frac{i}{\hbar}\left(\varepsilon_{\mathbf{p}-\mathbf{k}} - \varepsilon_{\mathbf{p}}\right)(t - t')\right] \times \exp\left[-i\zeta\left(\sin(\omega_0 t) - \sin(\omega_0 t')\right)\right], \tag{4}$$

where $\zeta = \mathbf{k}\cdot\mathbf{a}$, $\mathbf{a} = e\mathbf{E}_0/m\omega_0^2$ is the oscillation amplitude of the electrons driven by the RF (quiver amplitude), $\varepsilon_{\mathbf{p}} = p^2/2m$ is the electron energy with momentum $\mathbf{p}$. Here $f_{\mathbf{p}}$ is the equilibrium distribution function for the electron plasma.

Finally, using Eq. (2) and making a further Fourier transformation we obtain a solution for the potential φ in the form

$$\widetilde{\varphi}(\mathbf{k}, \omega) = \frac{4\pi\widetilde{\rho}_0(\mathbf{k}, \omega)}{k^2\varepsilon(k, \omega)}, \tag{5}$$

where we have introduced the frequency transforms $\widetilde{\varphi}(\mathbf{k}, \omega)$, $\widetilde{\rho}_0(\mathbf{k}, \omega)$ of the quantities

$$\begin{pmatrix} \widetilde{\rho}_0(\mathbf{k}, t) \\ \widetilde{\varphi}(\mathbf{k}, t) \end{pmatrix} = \begin{pmatrix} \rho_0(\mathbf{k}, t) \\ \varphi(\mathbf{k}, t) \end{pmatrix} e^{i\zeta\sin(\omega_0 t)}, \tag{6}$$

and $\varepsilon(k, \omega)$ is the RPA dielectric function (Lindhard, 1954; Lindhard & Winther, 1964).

In the case of a heavy particle with velocity $\mathbf{v}$ and charge Ze we neglect the effect of the RF on the particle and $\rho_0(\mathbf{r}, t) = Ze\delta(\mathbf{r} - \mathbf{v}t)$. We obtain

$$\widetilde{\rho}_0(\mathbf{k}, \omega) = 2\pi Ze\sum_{n=-\infty}^{\infty} J_n(\zeta)\delta\left(\omega - \mathbf{k}\cdot\mathbf{v} + n\omega_0\right), \tag{7}$$

where J_n is the Bessel function of nth order. Using Eqs. (5)–(7) for the self–consistent potential $\varphi(\mathbf{r}, t)$ we finally obtain

$$\varphi(\mathbf{r}, t) = \frac{Ze}{2\pi^2}\sum_{m,n=-\infty}^{\infty} e^{i(n-m)\omega_0 t}\int d\mathbf{k}\frac{e^{i\mathbf{k}\cdot(\mathbf{r}-\mathbf{v}t)}J_m(\zeta)J_n(\zeta)}{k^2\varepsilon(k, \mathbf{k}\cdot\mathbf{v} - n\omega_0)}. \tag{8}$$

This result represents the dynamical response of the medium to the motion of the test particle in the presence of the RF; it takes the form of an expansion over all the harmonics of the field frequency, with coefficients $J_n(\zeta)$ that depend on the intensity $W_L \propto a^2$.

From Eq. (8) it is straightforward to calculate the electric field $\mathbf{E}(\mathbf{r},t) = -\nabla\varphi(\mathbf{r},t)$, and the time average (with respect to the period $2\pi/\omega_0$ of the laser field) of the stopping field $\mathbf{E}_{\rm stop} = \langle \mathbf{E}(\mathbf{v}t,t)\rangle$ acting on the particle. Then, the averaged stopping power (SP) of the test particle becomes

$$S \equiv -Ze\frac{\mathbf{v}}{v}\cdot\mathbf{E}_{\rm stop} = \frac{2Z^2e^2}{(2\pi)^2 v}\sum_{n=-\infty}^{\infty}\int d\mathbf{k}\frac{\mathbf{k}\cdot\mathbf{v}}{k^2}J_n^2(\zeta)\,{\rm Im}\,\frac{-1}{\varepsilon(k,\Omega_n(\mathbf{k}))} \tag{9}$$

with $\Omega_n(\mathbf{k}) = n\omega_0 + \mathbf{k}\cdot\mathbf{v}$.

To illustrate the effects of the RF it is convenient to take into account the symmetry of the integrand in Eq. (9), with respect to the change $\mathbf{k}, n \to -\mathbf{k}, -n$. Using also the property of Bessel functions, $J_{-n}^2(\zeta) = J_n^2(\zeta)$, we obtain

$$S = \frac{Z^2e^2}{2\pi^2 v}\int d\mathbf{k}\frac{\mathbf{k}\cdot\mathbf{v}}{k^2}\left[J_0^2(\zeta)\,{\rm Im}\,\frac{-1}{\varepsilon(k,\mathbf{k}\cdot\mathbf{v})} + 2\sum_{n=1}^{\infty}J_n^2(\zeta)\,{\rm Im}\,\frac{-1}{\varepsilon(k,\Omega_n(\mathbf{k}))}\right]. \tag{10}$$

Hence, the SP depends on the particle velocity $\mathbf{v}$, the frequency ω_0 and the intensity $W_L = cE_0^2/8\pi$ of the RF (the intensity dependence is given through the quiver amplitude $\mathbf{a}$). Moreover, since the vector $\mathbf{k}$ in Eq. (10) is spherically integrated, S becomes also a function of the angle ϑ between the velocity $\mathbf{v}$, and the direction of polarization of RF, represented by $\mathbf{a}$.

By comparison, the SP in the absence of the RF is given by (Deutsch, 1986; Peter & Meyer-ter-Vehn, 1991)

$$S_B = \frac{Z^2e^2}{2\pi^2 v}\int d\mathbf{k}\frac{\mathbf{k}\cdot\mathbf{v}}{k^2}\,{\rm Im}\,\frac{-1}{\varepsilon\,(k,\mathbf{k}\cdot\mathbf{v})}. \tag{11}$$

In the presence of the RF the SP S_B is modified and is given by the first term in Eq. (10) (”no photon” SP)

$$S_0 = \frac{Z^2e^2}{2\pi^2 v}\int d\mathbf{k}\frac{\mathbf{k}\cdot\mathbf{v}}{k^2}J_0^2(\zeta)\,{\rm Im}\,\frac{-1}{\varepsilon\,(k,\mathbf{k}\cdot\mathbf{v})}. \tag{12}$$

Next we consider the case of a weak radiation field ($a < \lambda_s$, where λ_s is the characteristic screening length) at arbitrary angle ϑ between $\mathbf{v}$ and $\mathbf{E}_0$. In Eq. (10) we keep only the quadratic terms with respect to the quantity $\mathbf{a}$ and for the stopping power S we obtain

$$S = S_B + \frac{Z^2e^2}{4\pi^2 v}\int\frac{d\mathbf{k}}{k^2}(\mathbf{k}\cdot\mathbf{v})(\mathbf{k}\cdot\mathbf{a})^2\,{\rm Im}\left[\frac{1}{\varepsilon(k,\omega_0+\mathbf{k}\cdot\mathbf{v})} - \frac{1}{\varepsilon(k,\mathbf{k}\cdot\mathbf{v})}\right], \tag{13}$$

where S_B is the field-free SP given by Eq. (11). Note that due to the isotropy of the dielectric function $\varepsilon(k,\omega)$ the angular integrations in Eqs. (10)–(13) can be easily done.

It is well known that within classical description un upper cutoff parameter $k_{\max} = 1/r_{\min}$ (where $r_{\min}$ is the effective minimum impact parameter) must be introduced in Eqs. (11) and (13) to avoid the logarithmic divergence at large k. This divergence corresponds to the incapability of the classical perturbation theory to treat close encounters between the projectile particle and the plasma electrons properly. For $r_{\min}$ we use the effective minimum impact parameter excluding hard Coulomb collisions with a scattering angle larger than $\pi/2$. The resulting cutoff parameter $k_{\max} \simeq m(v^2+v_{\rm th}^2)/|Z|e^2$ is well known for energy loss calculations (see, e.g., Zwicknagel *et al.* (1999); Nersisyan *et al.* (2007) and references therein). Here $v_{\rm th}$ is the thermal velocity of the electrons. In particular, at low projectile velocities this cutoff parameter reads $k_{\max} = T/|Z|e^2$, where T is the plasma temperature given in energy units.

3 Energy loss of slow ions

In this section subsequent derivations are performed for the classical plasma and in the low–velocity limit of the ion. In this case the RPA dielectric function is given by (Fried & Conte, 1961)

$$\varepsilon(k,\omega) = 1 + \frac{1}{k^2\lambda_{\rm D}^2}W\left(\frac{\omega}{kv_{\rm th}}\right), \tag{14}$$

where $\lambda_{\rm D}$ is the Debye screening length, and $W(z) = g(z) + if(z)$ is the plasma dispersion function (Fried & Conte, 1961) with

$$g(z) = 1 - ze^{-z^2/2}\int_0^z e^{t^2/2}dt, \quad f(z) = \sqrt{\frac{\pi}{2}}ze^{-z^2/2}. \tag{15}$$

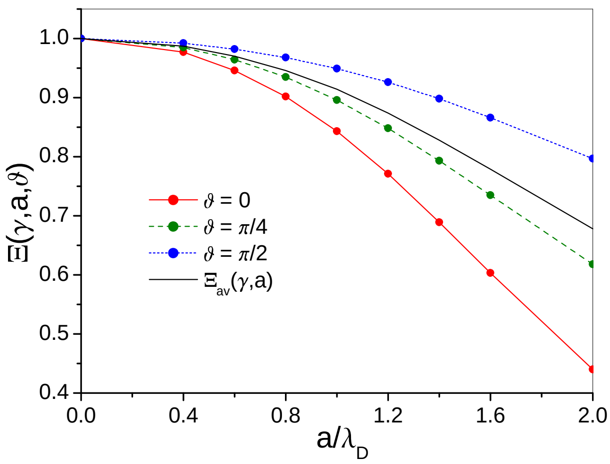

Figure 1: The dimensionless quantities $\Xi(\gamma,a,\vartheta)$ (the lines with symbols) and $\Xi_{\rm av}(\gamma,a)$ (the solid line without symbols) vs the intensity parameter of the laser field $a/\lambda_{\rm D}$ for $\vartheta=0$ (solid line), $\vartheta=\pi/4$ (dashed line), $\vartheta=\pi/2$ (dotted line) and for $\omega_0=1.2\omega_p$.

Consider now the SP determined by Eq. (10) in the limit of low–velocities, when $v\ll v_{\rm th}$. In this limit from Eqs. (10)–(15) we obtain

$$S(\gamma,a,\vartheta)=S_B\Xi(\gamma,a,\vartheta), \tag{16}$$

where

$$\Xi(\gamma,a,\vartheta)=\Xi_1(\gamma,a)+\Xi_2(\gamma,a)\sin^2\vartheta, \tag{17}$$

$$\Xi_s(\gamma,a)=\frac{6}{\psi(\xi)}\left\{\int_0^\xi\frac{k^3dk}{(k^2+1)^2}\int_0^1 J_0^2(Ak\mu)f_s(\mu)d\mu\right. \tag{18}$$

$$\left.+2\sqrt{\frac{2}{\pi}}\sum_{n=1}^{\infty}\int_0^\xi \mathrm{Im}\left[\frac{W_1(n/k\gamma)k^3dk}{(k^2+W(n/k\gamma))^2}\right]\int_0^1 J_n^2(Ak\mu)f_s(\mu)d\mu\right\}.$$

Here $s=1,2$, and $f_1(\mu)=\mu^2$, $f_2(\mu)=\frac{1}{2}(1-3\mu^2)$. Note that at the absence of the laser field (i.e., at $a\to 0$) $\Xi_1(\gamma,a)\to 1$, $\Xi_2(\gamma,a)\to 0$. In this case the SP is determined by the quantity S_B in Eq. (11) (Deutsch, 1986; Peter & Meyer-ter-Vehn, 1991)

$$S_B=\sqrt{\frac{2}{\pi}}\frac{Z^2e^2}{6\lambda_{\rm D}^2}\frac{v}{v_{\rm th}}\psi(\xi), \tag{19}$$

where

$$\psi(\xi)=\ln(1+\xi^2)-\frac{\xi^2}{1+\xi^2} \tag{20}$$

is the Coulomb logarithm with $\xi=k_{\max}\lambda_{\rm D}$. Also in Eqs. (16)–(18) we have introduced the angle ϑ between the velocity $\mathbf{v}$ and the polarization $\mathbf{a}$ vectors, $W_1(z)=dW(z)/dz$, $A=a/\lambda_D$, $\gamma=\omega_p/\omega_0<1$. Note that while the k integral in Eq. (11) diverges logarithmically in a field–free case, Eqs. (12) and (18) are finite and do not require any cutoff. The Bessel functions involved in these expressions due to the radiation field guarantee the convergence of the k–integrations. However, since in the sequel we shall compare Eqs. (16)–(18) with field–free SP S_B, for consistency the upper limits of the k–integrals in Eq. (18) are kept finite with the same upper cutoff parameter as in Eqs. (11) and (19).

In many experimental situations, the ions move in a plasma with random orientations of ϑ with respect to the direction of the polarization of laser field $\mathbf{a}$. The stopping power appropriate to this situation may be obtained by carrying out a spherical average over ϑ of $S(\gamma,a,\vartheta)$ in Eqs. (16) and (17). We find

$$S_{\rm av}(\gamma,a)=S_B\left[\Xi_1(\gamma,a)+\frac{2}{3}\Xi_2(\gamma,a)\right]\equiv S_B\Xi_{\rm av}(\gamma,a). \tag{21}$$

The study of the effect of a radiation field on the SP is easier in the case of low-intensities W_L when $a<\lambda_{\rm D}$. Then considering in Eqs. (16)–(18) only the quadratic terms with respect to a for the SP $S(\gamma,a,\vartheta)$ we obtain

$$S(\gamma,a,\vartheta)=S_B\left[1-\frac{a^2}{5\lambda_{\rm D}^2}(2\cos^2\vartheta+1)D(\gamma,\xi)\right], \tag{22}$$

where

$$D(\gamma,\xi)=\frac{1}{\psi(\xi)}\int_{1/\xi}^{\infty}\frac{dx}{x^3}\left\{\frac{1}{(x^2+1)^2}-\sqrt{\frac{2}{\pi}}\,\mathrm{Im}\left[\frac{W_1(x/\gamma)}{(1+x^2W(x/\gamma))^2}\right]\right\}. \tag{23}$$

Taking into account that $\gamma<1$ and $\xi\gg 1$ from Eqs. (22) and (23) we finally obtain $D(\gamma,\xi)\simeq 3/4\gamma^2$. It is seen that at low–velocities the SP $S(\gamma,a,\vartheta)$ decreases with the intensity of radiation field.

In Fig. 1 the quantities $\Xi(\gamma,a,\vartheta)$ and $\Xi_{\mathrm{av}}(\gamma,a)$ are shown vs the intensity parameter a/λ_{D} of the laser field for three values of angles $\vartheta=0$, $\vartheta=\pi/4$, $\vartheta=\pi/2$ and for $\omega_0=1.2\omega_p$. It is convenient to represent the intensity parameter a/λ_{D} in the form $a/\lambda_{\mathrm{D}}=0.18\lambda_0^2\sqrt{n_0W_L/T}$, where the wavelength ($\lambda_0$) and the intensity ($W_L$) of the laser field and the density (n_0) and the temperature (T) of plasma are measured in units μm, 10^{15} W/cm^2, 10^{20} cm^{-3} and keV, respectively. As an example consider the case when the electron quiver amplitude reaches the Debye screening length, $a=\lambda_{\mathrm{D}}$. For the values of the RF and plasma parameters with $\lambda_0=0.5\ \mu$m, $n_0=10^{18}$ cm^{-3}, $T=0.1$ keV, the above condition is fulfilled at the radiation field intensity $W_L=4.94\times10^{18}$ W/cm^2.

From Fig. 1 it is seen that the intense laser field may strongly reduce the SP of the low–velocity ion. And as expected the effect of the radiation field is maximal for $\vartheta=0$. Note that in this case and at $a=\lambda_{\mathrm{D}}$ the radiation field reduces the energy loss S_B approximately by 15 %. For explanation of the obtained result let us consider a simple physical model. The stopping power of the ion is defined as $S=-(1/v)\langle dW/dt\rangle$, where $\langle dW/dt\rangle$ is the averaged (with respect to the period of the radiation field) energy loss rate. We assume that the frequency of the radiation field ω_0 is larger than the effective frequency of the pairwise Coulomb collisions ν_{eff}. Also assuming that in the low–velocity limit the energy loss of the ion on the collective plasma excitations is negligible and is mainly determined by the Coulomb collisions we obtain $\langle dW/dt\rangle\sim\nu_{\mathrm{eff}}W$. On the other hand $\nu_{\mathrm{eff}}\sim 1/v_{\mathrm{eff}}^3$, where v_{eff} is the averaged relative velocity of the colliding particles. At $v<v_{\mathrm{th}}$ and for vanishing radiation field $v_{\mathrm{eff}}\simeq v_{\mathrm{th}}$. However, in the presence of the radiation field the averaged relative velocity of the collisions is $v_{\mathrm{eff}}\simeq(v_{\mathrm{th}}^2+v_E^2)^{1/2}$ and increases with the intensity of the laser field. Thus the effective collision frequency ν_{eff} and hence the stopping power of the ion are reduced with increasing the intensity of the radiation field.

At the end of this section we consider a practical example. Let us consider the stopping of the α–particles in the corona of the laser plasma. Although the thermonuclear reactions mainly occur far below the critical surface the stopping length of the α–particles is larger than the characteristic length scale of plasma inhomogeneity and some part of the α–particles transfer the energy to the plasma corona before they reach to the critical surface (Max, 1982). In the vicinity of the plasma critical density the intensity of the radiation field is very large and the stopping capacity of the plasma may be strongly reduced. In this example the typical temperature is $T=10$ keV and therefore $v_\alpha/v_{\mathrm{th}}=0.22$ ($E_\alpha=M_\alpha V_\alpha^2/2=3.5$ MeV, where E_α, M_α, v_α are the energy, the mass and the velocity of the α–particles). For $\lambda_0=0.5\ \mu$m, $W_L=2\times10^{17}$ W/cm^2, and $\omega_0=\omega_p\sqrt{2}$ (the plasma density is $n_0=n_c/2$, where n_c is the plasma critical density) we find $a\simeq\lambda_{\mathrm{D}}$. In this parameter regime the radiation field reduces the SP of the α–particles by 20 %.

4 Energy loss of fast ions

In this section we consider the energy loss of a fast heavy ion moving in a fully degenerate plasma (which means that the partially degenerate case could be postponed to a further presentation) in the presence of a radiation field. The longitudinal dielectric function of the degenerated electron gas is determined by Lindhard's expression (Lindhard, 1954; Lindhard & Winther, 1964). However, here we consider the simplest model of the dielectric function of a jellium. Previously a plasmon–pole approximation to $\varepsilon(k,\omega)$ for an electron gas was used for calculation of the SP (Basbas & Ritchie, 1982; Deutsch, 1995; Nersisyan & Das, 2000). In order to get easily obtainable analytical results, Basbas & Ritchie (1982) employed a simplified form that exhibits collective and single–particle effects

$$\mathrm{Im}\,\frac{-1}{\varepsilon(k,\omega)}=\pi\omega_p^2\frac{|\omega|}{\omega}\left[\delta\left(\omega^2-\omega_p^2\right)H(k_c-k)+\delta\left(\omega^2-\omega_k^2\right)H(k-k_c)\right], \tag{24}$$

where $H(x)$ is the Heaviside unit–step function, $\omega_k=\hbar k^2/2m$, $k_c=(2m\omega_p/\hbar)^{1/2}$, and ω_p is the plasma frequency. The cutoff parameter k_c is determined by equating the arguments of the two delta–functions in Eq. (24) at $k=k_c$. The first term in Eq. (24) describes the response due to nondispersive plasmon excitation in the region $k<k_c$, while the second term describes free–electron recoil in the range $k>k_c$ (single–particle excitations). Note that this approximate dielectric function satisfies at arbitrary k the usual frequency sum rule (Basbas & Ritchie, 1982; Deutsch, 1995; Nersisyan & Das, 2000).

In contrast to the previous section we consider here the fast projectile ion with $v\gtrsim v_c$ (where $v_c=\omega_p/k_c=(\hbar\omega_p/2m)^{1/2}$) which justifies the approximation (24) valid only in this specific case (Basbas & Ritchie, 1982).

It is constructive to consider first the case of a weak radiation field ($k_ca<1$) at arbitrary angle ϑ between $\mathbf{v}$ and $\mathbf{a}$. In this case the SP is determined by Eq. (13), where the field–free SP S_B in the high–velocity limit is given by (Lindhard, 1954; Lindhard & Winther, 1964; Deutsch, 1986, 1995)

$$S_B = \frac{Z^2 e^2 \omega_p^2}{v^2} \ln \left(\frac{2mv^2}{\hbar \omega_p} \right) . \tag{25}$$

Inserting Eq. (24) into (13) for the stopping power we obtain

$$S = \frac{2Z^2 \Sigma_0}{\lambda^2} \left\{ \ln \lambda + \frac{(k_c a)^2}{4} \left[\Phi_1(\lambda, \gamma) + \frac{1}{2} \Phi_2(\lambda, \gamma) \sin^2 \vartheta \right] \right\} , \tag{26}$$

where $\Sigma_0 = e^2 k_c^2 = 2\hbar\omega_p / a_0$, a_0 is the Bohr radius, $\Phi_1 = \Phi_{1c} + \Phi_{1s}$, $\Phi_2 = \Phi_{2c} + \Phi_{2s}$, $\lambda = v/v_c$, $\gamma = \omega_p/\omega_0 < 1$. Also

$$\Phi_{1c}(\lambda, \gamma) = \frac{1}{2\lambda^2} \left[\frac{6}{\gamma^2} \ln \lambda + \left(\frac{1}{\gamma} + 1 \right)^3 \ln \frac{\gamma}{1+\gamma} - \left(\frac{1}{\gamma} - 1 \right)^3 \ln \frac{\gamma}{1-\gamma} \right] , \tag{27}$$

$$\Phi_{2c}(\lambda, \gamma) = -3 \left[\Phi_{1c}(\lambda, \gamma) + \frac{1}{2\gamma^2 \lambda^2} \right] , \tag{28}$$

$$\Phi_{1s}(\lambda, \gamma) = \frac{1}{4\lambda^2} \left[\frac{1}{2} \left(\beta_1^2 + \eta_1^2 - \alpha_1^2 - \delta_1^2 \right) + \frac{3}{\gamma} \left(\beta_1 + \delta_1 - \alpha_1 - \eta_1 \right) \right. \tag{29}$$
$$\left. - \frac{1}{\gamma^3} \left(\frac{1}{\beta_1} - \frac{1}{\alpha_1} - \frac{1}{\eta_1} + \frac{1}{\delta_1} \right) + \frac{3}{\gamma^2} \ln \frac{\beta_1 \eta_1}{\alpha_1 \delta_1} + 1 - \lambda^4 \right] ,$$

$$\begin{aligned} \Phi_{2s}(\lambda, \gamma) = & \frac{\beta_1 - \alpha_1}{4} \left(1 - \frac{9}{\gamma \lambda^2} \right) + \frac{\eta_1 - \delta_1}{4} \left(1 + \frac{9}{\gamma \lambda^2} \right) \\ & - \frac{3}{8\lambda^2} \left(\beta_1^2 + \eta_1^2 - \alpha_1^2 - \delta_1^2 \right) + \frac{3}{4\gamma^3 \lambda^2} \left(\frac{1}{\beta_1} - \frac{1}{\alpha_1} - \frac{1}{\eta_1} + \frac{1}{\delta_1} \right) \\ & + \frac{1}{4\gamma} \left(\ln \frac{\beta_1 \delta_1}{\alpha_1 \eta_1} - \frac{9}{\gamma \lambda^2} \ln \frac{\beta_1 \eta_1}{\alpha_1 \delta_1} \right) + \frac{1}{4} \left(1 - \frac{1}{\lambda^2} \right) \left(\lambda^2 + 3 \right) , \end{aligned} \tag{30}$$

$$\begin{pmatrix} \alpha_n \\ \eta_n \end{pmatrix} = \max \left[\left(\frac{\lambda}{2} - \sqrt{\frac{\lambda^2}{4} \mp \frac{n}{\gamma}} \right)^2 ; \ 1 \right] , \tag{31}$$
$$\begin{pmatrix} \beta_n \\ \delta_n \end{pmatrix} = \left(\frac{\lambda}{2} + \sqrt{\frac{\lambda^2}{4} \mp \frac{n}{\gamma}} \right)^2 .$$

In Eq. (31) n is a positive integer ($n = 1, 2, ...$). The first term in Eq. (26) corresponds to the field–free SP (25) represented in a dimensionless form. The remaining terms proportional to the intensity of the radiation field (a^2), describe the collective (proportional to $\Phi_{1c;\,2c}(\lambda, \gamma)$) and single–particle (proportional to $\Phi_{1s;\,2s}(\lambda, \gamma)$) excitations. It should be noted that the stopping power Eq. (26) is not vanishing only at high–velocities when $\lambda \geqslant 2/\sqrt{\gamma}$.

Consider next the angular distribution of the SP at low–intensities of the RF. An analysis of the quantity $P = (S - S_B)/S_B$ (the relative deviation of S from S_B) for the proton projectile shows that at moderate velocities ($\lambda \gtrsim 2/\sqrt{\gamma}$) the angular distribution of P has a quadrupole nature. At $0 \leqslant \vartheta \leqslant \vartheta_0(\lambda, \gamma)$, where $\vartheta_0(\lambda, \gamma)$ is some value of the angle ϑ, the excitation of the waves with the frequencies $\omega_0 \pm \omega_p$ leads to the additional energy loss. At $\vartheta_0(\lambda, \gamma) \leqslant \vartheta \leqslant \pi/2$ the proton energy loss changes the sign and the total energy loss decreases. When the proton moves at the angle $\vartheta = \vartheta_0(\lambda, \gamma)$ with respect to the polarization vector $\mathbf{a}$ the radiation field has no any influence on the SP. However, at very large velocities ($\lambda \gg 2/\sqrt{\gamma}$) the relative deviation P is negative for arbitrary ϑ and the radiation field systematically reduces the energy loss of the proton.

Let us now investigate the influence of the intense radiation field on the stopping process when $\mathbf{v}$ is parallel to

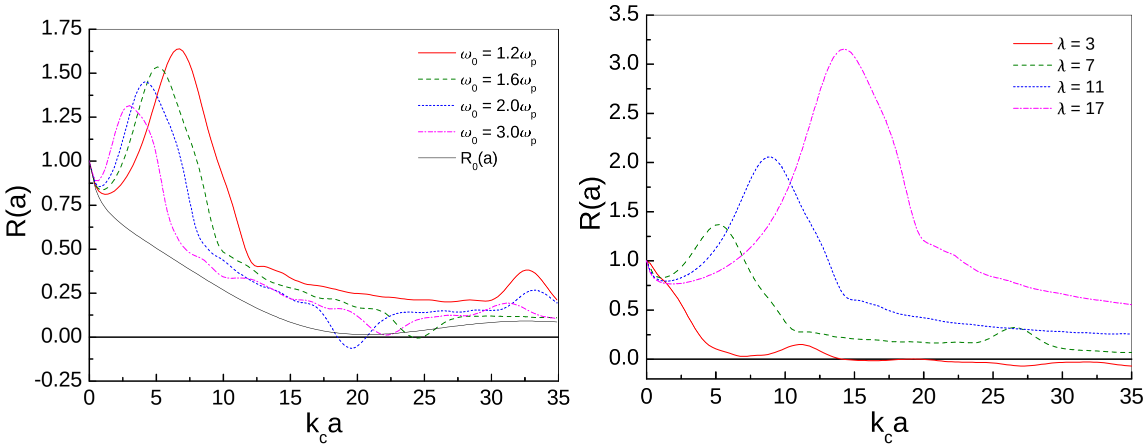

Figure 2: Left panel, the ratio $R(a) = S(a)/S_B$ as a function of dimensionless quantity $k_c a$ at $v = 8.6v_c$, $\omega_0 = 1.2\omega_p$ (solid line), $\omega_0 = 1.6\omega_p$ (dashed line), $\omega_0 = 2\omega_p$ (dotted line), $\omega_0 = 3\omega_p$ (dash–dotted line). Thin solid line corresponds to $R_0(a) = S_0(a)/S_B$ (see Eq. (33)). Right panel, same as in left panel but at $\omega_0 = 1.2\omega_p$, $v = 3v_c$ (solid line), $v = 7v_c$ (dashed line), $v = 11v_c$ (dotted line), $v = 17v_c$ (dash–dotted line).

a. It is expected that the effect of the RF is maximal in this case. From Eqs. (10) and (24) we obtain

$$S = S_0 + \frac{Z^2\Sigma_0}{\lambda^2}\left\{\sum_{n=1}^{n_-}\left(\frac{n}{\gamma}+1\right)J_n^2\left(Ap_n\right)\ln\frac{\lambda}{n/\gamma+1}\right.$$
$$-\sum_{n=1}^{n_+}\left(\frac{n}{\gamma}-1\right)J_n^2\left(Aq_n\right)\ln\frac{\lambda}{n/\gamma-1} \tag{32}$$
$$+\frac{1}{2}\sum_{n=1}^{N}\int_{\alpha_n(\lambda)}^{\beta_n(\lambda)}\frac{dx}{x^2}\left(\frac{n}{\gamma}+x\right)J_n^2\left(AP_n(x)\right)$$
$$\left.-\frac{1}{2}\sum_{n=1}^{\infty}\int_{\delta_n(\lambda)}^{\eta_n(\lambda)}\frac{dx}{x^2}\left(\frac{n}{\gamma}-x\right)J_n^2\left(AQ_n(x)\right)\right\},$$

where $A = k_c a$, $P_n(x) = (1/\lambda)(n/\gamma + x)$, $Q_n(x) = (1/\lambda)(n/\gamma - x)$, $p_n = P_n(1)$, $q_n = Q_n(1)$, and

$$S_0 = \frac{Z^2\Sigma_0}{\lambda^2}\left[J_0^2\left(\frac{A}{\lambda}\right)\ln\lambda + \frac{1}{2}\int_{1/\lambda}^{\lambda}\frac{dx}{x}J_0^2\left(Ax\right)\right] \tag{33}$$

is the SP without emission or absorption of the photons. Also we have introduced the notations

$$n_\pm = \text{int}\left(\frac{k_c v \pm \omega_p}{\omega_0}\right) = \text{int}\left[\gamma\left(\lambda \pm 1\right)\right], \tag{34}$$
$$N = \text{int}\left(\frac{mv^2}{2\hbar\omega_0}\right) = \text{int}\left(\frac{\gamma\lambda^2}{4}\right),$$

where $\text{int}(x)$ is the integer part of x. The quantities $\alpha_n(\lambda)$, $\beta_n(\lambda)$, $\delta_n(\lambda)$, $\eta_n(\lambda)$ in Eq. (32) are determined by Eq. (31). We note that in Eq. (32) the terms involving $n_\pm$ and N photons are not vanishing at $\lambda \geqslant 1/\gamma \mp 1$ and $\lambda \geqslant 2/\sqrt{\gamma}$, respectively. Similarly the SP (33) is not vanishing at $\lambda \geqslant 1$.

The first term in Eq. (33) describes the collective excitations while the second term corresponds to the single–particle excitations. From Eq. (33) it is seen that S_0 oscillates with the intensity of the laser field. However, the radiation field suppresses the excitation of the collective and the single–particle modes and the SP S_0 is less than the field–free SP S_B. As follows from Eq. (33) at high–intensities of the RF the SP S_0 is close to zero when $A/\lambda \simeq \mu_m$ (or alternatively at $\gamma(v_E/v) \simeq \mu_m$) with $m = 1, 2, \ldots$, where μ_m are the zeros of the Bessel function $J_0(\mu_m) = 0$ ($\mu_1 = 2.4$, $\mu_2 = 5.52$, $\mu_3 = 8.63\ldots$). Then the energy loss of the ion is mainly determined by the other terms in Eq. (32) and is stipulated by excitation of plasma waves with frequencies $n\omega_0 \pm \omega_p$. The first and the last pairs of terms in Eq. (32) describe the excitation of the collective and single–particle modes, respectively, with emission or absorption several photons. The number of photons ($n_\pm$, N) involved in the process of the inelastic interaction are determined by the energy–momentum conservations (see the arguments of the delta–functions in the dielectric function (24)).

The results of the numerical evaluation of the SP (Eqs. (32) and (33)) are shown in Fig. 2, where the ratio $R(a) = S(a)/S_B$ is plotted as a function of the laser field intensity ($k_c a = 5.38 W_L^{1/2} \omega_0^{-2} r_s^{-3/4}$, where r_s is the Wigner–Seitz density parameter and W_L and ω_0 are measured in units 10^{15} W/cm^2 and 10^{16} sec^{-1}, respectively). For instance, for Al target with $r_s = 2.07$, $\hbar\omega_p = 15.5$ eV, and $v_c = 1.2 \times 10^8$ cm/sec. From Fig. 2 it is seen that the SP exceeds the field–free SP and may change sign due to plasma irradiation by intense ($k_c a \gg 1$) laser field. Similar properties of the SP has been obtained previously for a classical plasma (Nersisyan & Akopyan, 1999). However, due to the higher density of the degenerate electrons (in metals typically $n_0 \sim 10^{23}$ cm^{-3}) the acceleration rate of the projectile particle is larger than similar rate in the case of a classical plasma. The acceleration effect occurs at $v_E/v \simeq \mu_m/\gamma$ (with $m = 1, 2, \dots$) when the SP S_0 nearly vanishes. It should be noted that in the laser irradiated plasma a parametrical instability is expected (Silin, 1973) with an increment increasing with the intensity of the radiation field. This restricts the possible acceleration time with stronger condition than in the case of a classical plasma. Finally, let us note that the effect of the enhancement of the SP of an ion moving in a laser irradiated plasma is intensified at smaller frequency (Fig. 2, left panel) of the radiation field ($\omega_0 \simeq \omega_p$ but $\omega_0 > \omega_p$) or at larger incident kinetic energy of the projectile ion (Fig. 2, right panel) when the numbers $n_\pm$ and N of the photons involved in the inelastic interaction process are strongly increased (Eq. (34)).

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