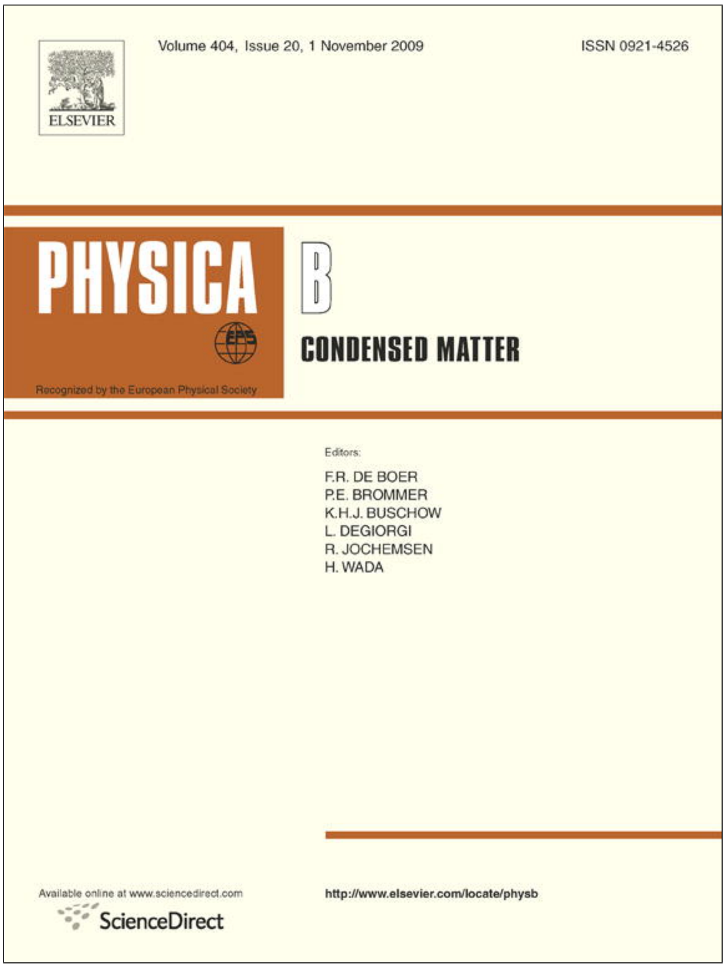

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Plasmons and polaritons in a semi-infinite plasma and a plasma slab

M. Apostol *, G. Vaman

Department of Theoretical Physics, Institute of Atomic Physics, Magurele-Bucharest Mg-6, PO Box Mg-35, Romania

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ABSTRACT

Plasmon and polariton modes are derived for an ideal semi-infinite (half-space) plasma and an ideal plasma slab by using a general, unifying procedure, based on equations of motion, Maxwell's equations and suitable boundary conditions. Known results are re-obtained in much a more direct manner and new ones are derived. The approach consists of representing the charge disturbances by a displacement field in the positions of the moving particles (electrons). The dielectric response and the electron energy loss are computed. The surface contribution to the energy loss exhibits an oscillatory behaviour in the transient regime near the surfaces. The propagation of an electromagnetic wave in these plasmas is treated by using the retarded electromagnetic potentials. The resulting integral equations are solved and the reflected and refracted waves are computed, as well as the reflection coefficient. For the slab we compute also the transmitted wave and the transmission coefficient. Generalized Fresnel's relations are thereby obtained for any incidence angle and polarization. Bulk and surface plasmon–polariton modes are identified. As it is well known, the field inside the plasma is either damped (evanescent) or propagating (transparency regime), and the reflection coefficient for a semi-infinite plasma exhibits an abrupt enhancement on passing from the propagating regime to the damped one (total reflection). Similarly, apart from characteristic oscillations, the reflection and transmission coefficients for a plasma slab exhibit an appreciable enhancement in the damped regime.

1. Introduction

After the discovery of bulk plasmons in an infinite electron plasma [1–3], there was a great deal of interest in plasmons occurring in structures with special geometries, like a half-space (semi-infinite) plasma, a plasma slab of finite thickness, a two-plasmas interface (two plasmas bounding each other), a two-dimensional sheet with an aperture, a slab with a cylindrical hole, structures with surface gratings or regular holes patterns, layered films, cylindrical rods and spherical particles, etc. There is a vast literature on various structures with special geometries exhibiting plasmon modes. These studies were aimed mainly at identifying new plasmon modes, like the surface plasmons [4–11], accounting for the electron energy loss experiments and exploring the interaction of the electron plasma with electromagnetic radiation (polariton excitations) [12–24]. More recently, a possible enhancement of the electromagnetic radiation scattered on electron plasmas with special geometries enjoyed a particular interest [25–27]. In all these studies the plasmon and polariton modes are of fundamental importance [28–32]. The methods used in deriving such results are of great diversity, resorting often to particular assumptions, such that the basic underlying mechanism of plasmons or polaritons' occurrence is often obscured. The need is therefore felt of having a general, unifying procedure for deriving plasmon and polariton modes in structures with special geometries, as based on the equation of motion of the charge density, Maxwell's equations and the corresponding boundary conditions. Such a procedure is presented in this paper for an ideal semi-infinite plasma and an ideal plasma slab.

We represent the charge disturbances as $\delta n = -n\,\mathrm{div}\,\mathbf{u}$, where n is the (constant, uniform) charge concentration and $\mathbf{u}$ is a displacement field of the mobile charges (electrons). This representation is valid for $\mathbf{Ku}(\mathbf{K}) \ll 1$, where $\mathbf{K}$ is the wavevector and $\mathbf{u}(\mathbf{K})$ is the Fourier component of the displacement field. We assume a rigid neutralizing background of positive charge, as in the well-known jellium model. In the static limit, i.e. for Coulomb interaction, the Lagrangian of the electrons can be written as

$$L = \int d\mathbf{r}\left[\frac{1}{2}mn\dot{\mathbf{u}}^2 - \frac{1}{2}\int d\mathbf{r}' U(|\mathbf{r}-\mathbf{r}'|)\delta n(\mathbf{r})\delta n(\mathbf{r}')\right] + e\int d\mathbf{r}\Phi(\mathbf{r})\delta n(\mathbf{r}), \tag{1}$$

where m is the electron mass, $U(r) = e^2/r$ the Coulomb energy, $-e$ the electron charge and $\Phi(\mathbf{r})$ the external scalar potential. Eq. (1)

* Corresponding author.
E-mail address: apoma@theory.nipne.ro (M. Apostol).

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leads to the equation of motion

$$m\ddot{\mathbf{u}} = n\,\mathrm{grad}\int d\mathbf{r}' U(|\mathbf{r}-\mathbf{r}'|)\,\mathrm{div}\,\mathbf{u}(\mathbf{r}') + e\,\mathrm{grad}\,\Phi \tag{2}$$

which is the starting equation of our approach. We leave aside the dissipation effects (which can easily be included in Eq. (2)).

By using the Fourier transform for an infinite plasma it is easy to see that the eigenmode of the homogeneous Eq. (2) is the well-known bulk plasmon mode given by $\omega_p^2 = 4\pi ne^2/m$. On the other side, equation $\delta n = -n\,\mathrm{div}\,\mathbf{u}$ is equivalent with Maxwell's equation $\mathrm{div}\,\mathbf{E}_i = -4\pi e\delta n$, where $\mathbf{E}_i = 4\pi ne\mathbf{u}$ is the internal electric field (equal to $-4\pi\mathbf{P}$, where $\mathbf{P}$ is the polarization). Making use of the electric displacement $\mathbf{D} = -\mathrm{grad}\,\Phi = \varepsilon(\mathbf{D}+\mathbf{E}_i)$, we get the well-known dielectric function $\varepsilon = 1-\omega_p^2/\omega^2$ in the long-wavelength limit from the solution of the inhomogeneous Eq. (2). Similarly, since the current density is $\mathbf{j} = -en\dot{\mathbf{u}}$, we get the well-known electrical conductivity $\sigma = i\omega_p^2/4\pi\omega$.

We apply this approach to a semi-infinite plasma and a plasma slab. First, we derive the surface and bulk plasmon modes and obtain the dielectric response and the electron energy loss for a semi-infinite plasma. The surface contribution to the energy loss exhibits an oscillatory behaviour in the transient regime near the surface. Further on, we consider the interaction of the semi-infinite plasma with the electromagnetic field, as described by the usual term $(1/c)\int d\mathbf{r}\mathbf{j}\mathbf{A} - \int d\mathbf{r}\rho\Phi$ in the Lagrangian, where $\mathbf{A}$ is the vector potential, $\rho = en\,\mathrm{div}\,\mathbf{u}$ is the charge density and Φ is the scalar potential. We limit ourselves to the interaction with the electric field, and compute the reflected and refracted waves, as well as the reflection coefficient. Generalized Fresnel's relations are obtained for any incidence angle and polarization. We find it more convenient to use the radiation formulae for the retarded potentials, instead of using directly the Maxwell's equations, and the resulting integral equations are solved. Bulk and surface plasmon–polariton modes are identified. The field inside the plasma is either damped (evanescent) or propagating (transparency regime), and the reflection coefficient exhibits an abrupt enhancement on passing from the propagating to the damping regime (total reflection). Finally, we give similar results for a plasma slab, where we compute also the transmitted field and the transmission coefficient. Apart from characteristic oscillations, the reflection and transmission coefficients for a plasma slab exhibit an appreciable enhancement in the damped regime. The present approach can be extended to various other plasma structures with special geometries.

2. Plasma eigenmodes

We consider an ideal semi-infinite plasma extending over the half-space $z>0$ (and bounded by the vacuum for $z<0$). The displacement field $\mathbf{u}$ is then represented as $(\mathbf{v}, u_3)\theta(z)$, where $\mathbf{v}$ is the displacement component in the (x,y)-plane, u_3 is the displacement component along the z-direction and $\theta(z)=1$ for $z>0$ and $\theta(z)=0$ for $z<0$ is the step function. In equation of motion (2) $\mathrm{div}\,\mathbf{u}$ is then replaced by

$$\mathrm{div}\,\mathbf{u} = \left(\mathrm{div}\,\mathbf{v} + \frac{\partial u_3}{\partial z}\right)\theta(z) + u_3(0)\delta(z), \tag{3}$$

where $u_3(0) = u_3(\mathbf{r}, z=0)$, $\mathbf{r}$ being the in-plane (x,y) position vector. Eq. (2) becomes

$$m\ddot{\mathbf{u}} = ne^2\,\mathrm{grad}\int d\mathbf{r}'dz' \frac{1}{\sqrt{(\mathbf{r}-\mathbf{r}')^2+(z-z')^2}} \times\left[\mathrm{div}\,\mathbf{v}(\mathbf{r}'.z') + \frac{\partial u_3(\mathbf{r}',z')}{\partial z'}\right]$$
$$+ne^2\,\mathrm{grad}\int d\mathbf{r}' \frac{1}{\sqrt{(\mathbf{r}-\mathbf{r}')^2+z^2}} u_3(\mathbf{r}',0) + e\,\mathrm{grad}\,\Phi \tag{4}$$

for $z>0$. One can see the (de)-polarizing field occurring at the free surface $z=0$ (the second integral in Eq. (4)).

We use Fourier transforms of the type

$$\mathbf{u}(r,z;t) = \sum_{\mathbf{k}}\int d\omega\mathbf{u}(\mathbf{k},z;\omega)e^{i\mathbf{k}\mathbf{r}}e^{-i\omega t} \tag{5}$$

(for in-plane unit area), as well as the Fourier representation

$$\frac{1}{\sqrt{r^2+z^2}} = \sum_{\mathbf{k}}\frac{2\pi}{k}e^{-k|z|}e^{i\mathbf{k}\mathbf{r}} \tag{6}$$

for the Coulomb potential. Then, it is easy to see that Eq. (4) leads to the integral equation

$$\omega^2 v = \frac{1}{2}k\omega_p^2\int_0^\infty dz' v e^{-k|z-z'|} + \frac{1}{2k}\omega_p^2\int_0^\infty dz' \frac{\partial v}{\partial z'}\frac{\partial}{\partial z'}e^{-k|z-z'|} - \frac{iek}{m}\Phi \tag{7}$$

and $iku_3 = \partial v/\partial z$, where we have dropped out for simplicity the arguments $\mathbf{k}, z$ and ω. The $\mathbf{v}$-component of the displacement field is directed along the wavevector $\mathbf{k}$ (in-plane longitudinal waves). This integral equation can easily be solved. Integrating by parts in its *rhs* we get

$$\omega^2 v = \omega_p^2 v - \frac{1}{2}\omega_p^2 v_0 e^{-kz} - \frac{iek}{m}\Phi, \tag{8}$$

hence

$$v = \frac{iek\omega_p^2}{m}\frac{\Phi_0}{(\omega^2-\omega_p^2)(2\omega^2-\omega_p^2)}e^{-kz} - \frac{iek}{m}\frac{\Phi}{\omega^2-\omega_p^2}$$
$$u_3 = -\frac{ek\omega_p^2}{m}\frac{\Phi_0}{(\omega^2-\omega_p^2)(2\omega^2-\omega_p^2)}e^{-kz} - \frac{e}{m}\frac{\Phi'}{\omega^2-\omega_p^2}, \tag{9}$$

where $v_0 = v(z=0)$, $\Phi_0 = \Phi(z=0)$ and $\Phi' = \partial\Phi/\partial z$. One can see the surface contributions (terms proportional to $\Phi_0 e^{-kz}$) and bulk contributions (Φ, Φ'-terms).

The solutions given by Eqs. (9) exhibit two eigenmodes, the bulk plasmon $\omega_b = \omega_p$ and the surface plasmon $\omega_s = \omega_p/\sqrt{2}$, as it is well known. Indeed, the homogeneous Eq. (8) ($\Phi = 0$) has two solutions: the surface plasmon $v = v_0 e^{-kz}$ for $\omega^2 = \omega_p^2/2$ and the bulk plasmon $v_0 = 0$ for $\omega^2 = \omega_p^2$. Making use of this observation we can represent the general solution as an eigenmodes series

$$v(\mathbf{k},z) = \sqrt{2k}v_0(\mathbf{k})e^{-kz} + \sum_\kappa\sqrt{\frac{2k^2}{\kappa^2+k^2}}v(\mathbf{k},\kappa)\sin\kappa z \tag{10}$$

for $z>0$, where $v(\mathbf{k},-\kappa) = -v(\mathbf{k},\kappa)$ and $iku_3(\mathbf{k},z) = \partial v(\mathbf{k},z)/\partial z$. Then, it is easy to see that the hamiltonian $H = T+U$ corresponding to the Lagrangian $L = T-U$ given by Eq. (1) becomes

$$T = nm\sum_{\mathbf{k}}\dot{v}_0^*(\mathbf{k})\dot{v}_0(\mathbf{k}) + nm\sum_{\mathbf{k}\kappa}\dot{v}^*(\mathbf{k},\kappa)\dot{v}(\mathbf{k},\kappa)$$
$$U = 2\pi n^2e^2\sum_{\mathbf{k}}v_0^*(\mathbf{k})v_0(\mathbf{k}) + 4\pi n^2e^2\sum_{\mathbf{k}\kappa}v^*(\mathbf{k},\kappa)v(\mathbf{k},\kappa), \tag{11}$$

where T is the kinetic energy and U is the potential energy. We can see that this hamiltonian corresponds to harmonic oscillators with frequencies $\omega_s = \omega_p/\sqrt{2}$ and $\omega_b = \omega_p$.

Making use of $\mathbf{E}_i = 4\pi ne\mathbf{u}$ and Eqs. (9) we can write down the internal field (polarization) as

$$E_\perp(\mathbf{k},z;\omega) = \frac{ik\omega_p^4\Phi(\mathbf{k},0;\omega)}{(\omega^2-\omega_p^2)(2\omega^2-\omega_p^2)}e^{-kz} - \frac{ik\omega_p^2\Phi(\mathbf{k},z;\omega)}{\omega^2-\omega_p^2}$$
$$E_\parallel(\mathbf{k},z;\omega) = -\frac{k\omega_p^4\Phi(\mathbf{k},0;\omega)}{(\omega^2-\omega_p^2)(2\omega^2-\omega_p^2)}e^{-kz} - \frac{\omega_p^2\Phi'(\mathbf{k},z;\omega)}{\omega^2-\omega_p^2}, \tag{12}$$

where $E_\perp$ is directed along the in-plane wavevector $\mathbf{k}$ and $E_\parallel$ is parallel with the z-axis (perpendicular to the surface $z=0$). This is the dielectric response of the semi-infinite plasma to an external potential.

We take an external potential of the form $\Phi(\mathbf{k},z)=\Phi^0(\mathbf{k})e^{i\kappa z}$ (leaving aside the frequency argument ω) and get the electric displacement $\mathbf{D}_\perp(\mathbf{k},z)=-i\mathbf{k}\Phi^0(\mathbf{k})e^{i\kappa z}$ and $D_\parallel(\mathbf{k},z)=-i\kappa\Phi^0(\mathbf{k})e^{i\kappa z}$ from $\mathbf{D}=-\mathrm{grad}\,\Phi$. We can see that the surface terms do not contribute to this response, as expected, since these terms are localized. Making use of $\mathbf{E}_i=(1/\varepsilon-1)\mathbf{D}$, we get the well-known dielectric function $\varepsilon(\kappa,\omega)=1-\omega_p^2/\omega^2$ in the long-wavelength limit.

3. Electron energy loss

It is well known that the energy loss per unit time (stopping power) is given by

$$P=\frac{d}{dt}\left(\frac{mv^2}{2}\right)=-e\mathbf{v}\mathbf{E}_i \tag{13}$$

for an electron moving with velocity $\mathbf{v}=(\mathbf{v}_\perp,v_\parallel)$, where the field $\mathbf{E}_i$ is taken at $\mathbf{r}=\mathbf{v}_\perp t$ and $z=v_\parallel t$ for $t>0$ ($z>0$). It is assumed that the electron energy is sufficiently large and the energy loss is small enough to use a constant $\mathbf{v}$ in estimating the *rhs* of Eq. (13). The potential created by the electron is given by the Poisson equation $\Delta\Phi=4\pi e\delta(\mathbf{r}-\mathbf{v}_\perp t)\delta(z-v_\parallel t)$, whence, by making use of the Fourier representation (6), we get

$$\Phi(\mathbf{k},z;\omega)=-\frac{2ev_\parallel}{(\omega-\mathbf{k}\mathbf{v}_\perp)^2+k^2v_\parallel^2}e^{-i(\mathbf{k}\mathbf{v}_\perp-\omega)z/v_\parallel}. \tag{14}$$

We introduce this potential in Eqs. (12) and compute the energy loss given by Eq. (13). It contains two contributions, one associated with the bulk plasmons,

$$P_b=e^2\omega_p^2\sum_{\mathbf{k}}\int d\omega\frac{i\omega}{\omega_p^2-\omega^2}\cdot\frac{2v_\parallel}{(\omega-\mathbf{k}\mathbf{v}_\perp)^2+k^2v_\parallel^2} \tag{15}$$

and another arising from surface effects,

$$P_s=e^2\omega_p^4\sum_{\mathbf{k}}\int d\omega\frac{1}{(\omega^2-\omega_p^2/2)(\omega^2-\omega_p^2)}\cdot\frac{v_\parallel(i\mathbf{k}\mathbf{v}_\perp-kv_\parallel)}{(\omega-\mathbf{k}\mathbf{v}_\perp)^2+k^2v_\parallel^2}\times e^{-kv_\parallel t}e^{i(\mathbf{k}\mathbf{v}_\perp-\omega)t}. \tag{16}$$

In performing the ω-integrations in Eqs. (15) and (16) we retain only the plasmon contributions arising from the poles $\omega=\omega_p$ and $\omega=\omega_p/\sqrt{2}$. For normal incidence ($v_\perp=0$, $v_\parallel=v$) we get easily the well-known bulk contribution $P_b=(-e^2\omega_p^2/v)\ln(vk_0/\omega_p)$, where k_0 is an upper cut-off (associated, as usually, with the ionization energy, or with the inverse of the mean inter-particle spacing, etc), and the surface contribution

$$P_s=-\frac{e^2\omega_p}{vt}(\sqrt{2}\sin\omega_p t/\sqrt{2}-\sin\omega_p). \tag{17}$$

We can see in Eq. (17) the oscillatory behaviour of the stopping power arising from the surface effects in the transient regime near the surface.

4. Interaction with the electromagnetic field. Polaritons

We assume a plane wave incident on the plasma surface under angle α. Its frequency is given by $\omega=cK$, where c is the velocity of light and the wavevector $\mathbf{K}=(\mathbf{k},\kappa)$ has the in-plane component $\mathbf{k}$ and the perpendicular-to-plane component κ, such as $k=K\sin\alpha$ and $\kappa=K\cos\alpha$. In addition, $\mathbf{k}=k(\cos\varphi,\sin\varphi)$. The electric field is taken as $\mathbf{E}_0=E_0(\cos\beta,0,-\sin\beta)\times e^{i\mathbf{kr}}e^{i\kappa z}e^{-i\omega t}$, and we impose the condition $\cos\beta\sin\alpha\cos\varphi-\sin\beta\cos\alpha=0$ (transversality condition $\mathbf{K}\mathbf{E}_0=0$). The angle β defines the direction of the polarization of the incident field. The geometry of the incident wave is shown in Fig. 1.

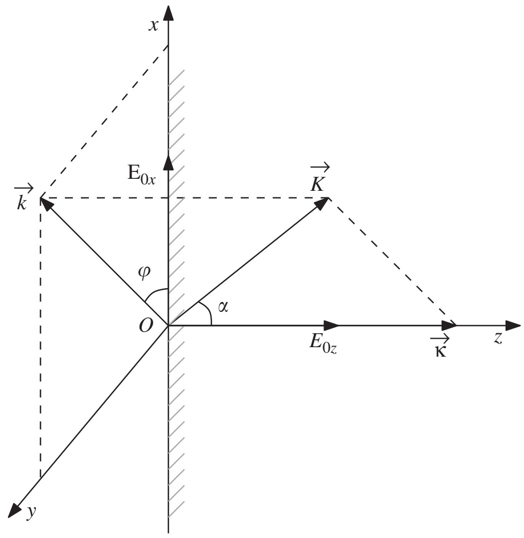

Fig. 1. Electromagnetic plane wave $\mathbf{E}_0$, with wavevector $\mathbf{K}$, incident on the surface $z=0$.

In the presence of an electromagnetic wave we use the equation of motion

$$\omega^2\mathbf{u}=\frac{e}{m}\mathbf{E}+\frac{e}{m}\mathbf{E}_0e^{i\kappa z} \tag{18}$$

for $z>0$, where $\mathbf{E}$ is the polarizing field; in Eq. (18) we have preserved explicitly only the z-dependence (i.e. we leave aside the factors $e^{i\mathbf{kr}}e^{-i\omega t}$). We find it convenient to employ the vector potential

$$\mathbf{A}(\mathbf{r},z;t)=\frac{1}{c}\int d\mathbf{r}'\int dz'\frac{\mathbf{j}(\mathbf{r}',z';t-R/c)}{R} \tag{19}$$

and the scalar potential

$$\Phi(\mathbf{r},z;t)=\int d\mathbf{r}'\int dz'\frac{\rho(\mathbf{r}',z';t-R/c)}{R}, \tag{20}$$

where $\mathbf{j}=-ne\dot{\mathbf{u}}\theta(z)e^{i\mathbf{kr}}e^{-i\omega t}$ is the current density, $\rho=ne\,\mathrm{div}\,\mathbf{u}=ne(i\mathbf{k}\mathbf{v}+(\partial u_3/\partial z))\theta(z)e^{i\mathbf{kr}}e^{-i\omega t}+neu_3(0)\delta(z)e^{i\mathbf{kr}}e^{-i\omega t}$ is the charge density and $R=\sqrt{(\mathbf{r}-\mathbf{r}')^2+(z-z')^2}$. The integrals in Eqs. (19) and (20) implies the known integral [33]

$$\int_{|z|}^{\infty}dxJ_0(k\sqrt{x^2-z^2})e^{i\omega x/c}=\frac{i}{\kappa}e^{i\kappa|z|}, \tag{21}$$

where J_0 is the zeroth-order Bessel function of the first kind (and $\omega^2/c^2=\kappa^2+k^2$). It is convenient to use the projections of the in-plane displacement field $\mathbf{v}$ on the vectors $\mathbf{k}$ and $\mathbf{k}_\perp=k(-\sin\varphi,\cos\varphi)$, $\mathbf{k}_\perp\mathbf{k}=0$. We denote these components by $v_1=\mathbf{k}\mathbf{v}/k$ and $v_2=\mathbf{k}_\perp\mathbf{v}/k$, and use also the components $E_1=\mathbf{k}\mathbf{E}/k$, $E_2=\mathbf{k}_\perp\mathbf{E}/k$ and similar ones for the external field $\mathbf{E}_0$. We give here the components of the external field

$$E_{01}=E_0\cos\beta\cos\varphi,\quad E_{02}=-E_0\cos\beta\sin\varphi,\quad E_{03}=-E_0\sin\beta. \tag{22}$$

One can check immediately the transversality condition $E_{01}k+E_{03}\kappa=0$. Making use of $\mathbf{E}=-(1/c)(\partial\mathbf{A}/\partial t)-\mathrm{grad}\,\Phi$, Eqs. (19) and (20) give the electric field

$$E_1=-2\pi ine\kappa\int_0 dz'v_1(z')e^{i\kappa|z-z'|}-2\pi ne\frac{k}{\kappa}\int_0 dz'u_3(z')\frac{\partial}{\partial z'}e^{i\kappa|z-z'|},$$

$$E_2 = -2\pi ine\frac{\omega^2}{c^2\kappa}\int_0 dz' v_2(z')e^{i\kappa|z-z'|},$$

$$E_3 = 2\pi ne\frac{k}{\kappa}\int_0 dz' v_1(z')\frac{\partial}{\partial z}e^{i\kappa|z-z'|} - 2\pi ine\frac{k^2}{\kappa}\int_0 dz' u_3(z')e^{i\kappa|z-z'|} + 4\pi neu_3 \tag{23}$$

for $z>0$. It is worth observing in deriving these equations the non-intervertibility of the derivatives and the integrals, according to the identity

$$\frac{\partial}{\partial z}\int_0 dz' f(z')\frac{\partial}{\partial z'}e^{i\kappa|z-z'|} = \kappa^2\int_0 dz' f(z')e^{i\kappa|z-z'|} - 2i\kappa f(z) \tag{24}$$

for any function $f(z)$, $z>0$; it is due to the discontinuity in the derivative of the function $e^{i\kappa|z-z'|}$ for $z=z'$. Now, we employ equation of motion (18) in Eqs. (23) and get the integral equations

$$\omega^2 v_1 = -\frac{i\omega_p^2\kappa}{2}\int_0 dz' v_1(z')e^{i\kappa|z-z'|} - \frac{\omega_p^2 k}{2\kappa}\int_0 dz' u_3(z')\frac{\partial}{\partial z'}e^{i\kappa|z-z'|} + \frac{e}{m}E_{01}e^{i\kappa z},$$

$$\omega^2 v_2 = -\frac{i\omega_p^2\omega^2}{2c^2\kappa}\int_0 dz' v_2(z')e^{i\kappa|z-z'|} + \frac{e}{m}E_{02}e^{i\kappa z},$$

$$\omega^2 u_3 = \frac{\omega_p^2 k}{2\kappa}\int_0 dz' v_1(z')\frac{\partial}{\partial z}e^{i\kappa|z-z'|} - \frac{i\omega_p^2 k^2}{2\kappa}\int_0 dz' u_3(z')e^{i\kappa|z-z'|} + \omega_p^2 u_3 + \frac{e}{m}E_{03}e^{i\kappa z} \tag{25}$$

for the coordinates $v_{1,2}$ and u_3 in the region $z>0$.

The second Eq. (25) can be solved straightforwardly by noticing that

$$\frac{\partial^2}{\partial z^2}\int_0 dz' v_2(z')e^{i\kappa|z-z'|} = -\kappa^2\int_0 dz' v_2(z')e^{i\kappa|z-z'|} + 2i\kappa v_2. \tag{26}$$

We get

$$\frac{\partial^2 v_2}{\partial z^2} + (\kappa^2 - \omega_p^2/c^2)v_2 = 0. \tag{27}$$

The solution of this equation is

$$v_2 = \frac{2eE_{02}}{m\omega_p^2}\cdot\frac{\kappa(\kappa-\kappa')}{K^2}e^{i\kappa' z}, \tag{28}$$

where

$$\kappa' = \sqrt{\kappa^2 - \omega_p^2/c^2} = \frac{1}{c}\sqrt{\omega^2\cos^2\alpha - \omega_p^2}. \tag{29}$$

The wavevector κ' can also be written in a more familiar form $\kappa' = (\omega/c)\sqrt{\varepsilon - \sin^2\alpha}$, where $\varepsilon = 1 - \omega_p^2/\omega^2$ is the dielectric function. The corresponding component of the (total) electric field (the refracted field), can be obtained from Eq. (18); it is given by $(m\omega^2/e)v_2$. For $\kappa^2 < \omega_p^2/c^2$ ($\omega\cos\alpha < \omega_p$) this field does not propagate. For $\kappa^2 > \omega_p^2/c^2$ (ω greater than the transparency edge $\omega_p/\cos\alpha$) it represents a refracted wave (transparency regime) with the refraction angle α' given by Snell's law

$$\frac{\sin\alpha'}{\sin\alpha} = \frac{1}{\sqrt{1-\omega_p^2/\omega^2}} = 1/\sqrt{\varepsilon}. \tag{30}$$

The polariton frequency is given by

$$\omega^2 = c^2K^2 = \omega_p^2 + c^2K'^2 \tag{31}$$

as it is well known, where $K'^2 = \kappa'^2 + k^2$.

The first and the third Eqs. (25) can be solved by using an equation similar with Eq. (26) and by noticing that they imply

$$\kappa'^2 u_3 = ik\frac{\partial v_1}{\partial z}. \tag{32}$$

We get

$$v_1 = \frac{2eE_{01}}{m\omega_p^2}\cdot\frac{\kappa'(\kappa-\kappa')}{\kappa\kappa'+k^2}e^{i\kappa' z} \tag{33}$$

and

$$u_3 = \frac{2eE_{03}}{m\omega_p^2}\cdot\frac{\kappa(\kappa-\kappa')}{\kappa\kappa'+k^2}e^{i\kappa' z}. \tag{34}$$

Similarly, the corresponding components of the refracted field are given by Eq. (18). It is easy to check the transversality condition $v_1k + u_3\kappa' = 0$ (and the vanishing of the bulk charge $ne(i\mathbf{kv} + \partial u_3/\partial z) = 0$).

We can see that the polarization field $\mathbf{E}$ in Eq. (18) cancels out the original incident field $\mathbf{E}_0$ and gives the total, refracted field $m\omega^2\mathbf{u}/e$ inside the plasma. This is an illustration of the so-called Ewald–Oseen extinction theorem [17,34].

It is worth investigating the eigenvalues of the homogeneous system of integral Eqs. (25), for parameter κ given by $\kappa = \sqrt{\omega^2/c^2 - k^2}$. Such eigenvalues are given by the roots of the vanishing denominator in Eqs. (33) and (34), i.e. by equation $\kappa\kappa' + k^2 = 0$. This equation has real roots for ω only for the damped regime, i.e. for $\kappa = i|\kappa|$ and $\kappa' = i|\kappa'|$. Providing these conditions are satisfied, there is only one acceptable branch of excitations, given by

$$\omega^2 = \frac{2\omega_p^2c^2k^2}{\omega_p^2 + 2c^2k^2 + \sqrt{\omega_p^4 + 4c^4k^4}}. \tag{35}$$

We can see that $\omega\sim ck$ in the long wavelength limit and it approaches the surface-plasmon frequency $\omega\sim\omega_p/\sqrt{2}$ in the non-retarded limit ($ck\to\infty$). These excitations are surface plasmon–-polariton modes. We note that they imply $v_2 = 0$ and $v_1, u_3 \sim e^{-|\kappa'|z}$. In addition, a careful analysis of the homogeneous system of Eqs. (25) reveals another branch of excitations, given by $\omega = \omega_p$, which, occurring in this context, may be termed the bulk plasmon–polariton modes. They are characterized by $v_2 = 0$ and $v_1(\mathbf{k}, 0) = 0$. For all these modes we have $u_3 = [ic^2k/(\omega^2 - c^2k^2 - \omega_p^2)]\partial v_1/\partial z$.

In order to get the reflected wave (the region $z<0$) we turn to Eqs. (23) and use therein the solutions given above for $v_{1,2}$ and u_3. It is worth noting here that the discontinuity term $\omega_p^2u_3$ does not appear anymore in these equations (because $z'>0$ and $z<0$ and we cannot have $z=z'$). The integrations in Eqs. (23) are straightforward and we get the field

$$E_1 = E_{01}\frac{\kappa-\kappa'}{\kappa+\kappa'}\cdot\frac{\kappa\kappa'-k^2}{\kappa\kappa'+k^2}e^{-i\kappa z}, \tag{36}$$

$$E_2 = E_{02}\frac{\kappa-\kappa'}{\kappa+\kappa'}e^{-i\kappa z} \tag{37}$$

and

$$E_3 = -E_{03}\frac{\kappa-\kappa'}{\kappa+\kappa'}\cdot\frac{\kappa\kappa'-k^2}{\kappa\kappa'+k^2}e^{-i\kappa z}. \tag{38}$$

We can see that this field represents the reflected wave ($\kappa\to-\kappa$) and we can check its transversality to the propagation wavevector. Making use of the reflected field $\mathbf{E}_{refl}$ given by Eqs. (36)–(38) and the refracted field $\mathbf{E}_{refr}$ obtained from Eqs. (18) and (23) ($\mathbf{E}_{refr} = \mathbf{E} + \mathbf{E}_0 = m\omega^2\mathbf{u}/e$) one can check the continuity of the electric field and electric displacement at the surface ($z=0$) in the form $E_{1,2refl} + E_{01,2} = E_{1,2refr}$, $E_{3refl} + E_{03} = \varepsilon E_{3refr}$, where $\varepsilon = 1 - \omega_p^2/\omega^2$. The angle of total polarization (Brewster's angle) is given by $\kappa\kappa' - k^2 = 0$, or $\tan^2\alpha = 1 - \omega_p^2/\omega^2 = \varepsilon$ (for $\alpha < \pi/4$). The above equations provide generalized Fresnel's relations between the amplitudes of the reflected, refracted and incident waves at the surface for any incidence angle and polarization. They can also be

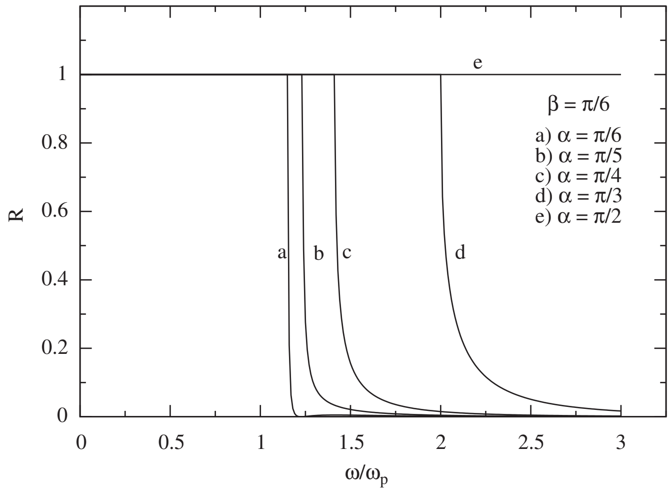

Fig. 2. Reflection coefficient for a semi-infinite plasma for $\beta=\pi/6$ and various incidence angles α. One can see the shoulder occurring at the transparency edge $\omega_p/\cos\alpha$ and the zero occurring at $\omega^2=\omega_p^2/(1-\tan^2\alpha)$ for $\alpha=\beta=\pi/6$ ($R_2=0,\varphi=0$).

written by using $\omega^2=\omega_p^2/(1-\varepsilon)$, where ε is the dielectric function.

The reflection coefficient $R=|\mathbf{E}_{refl}|^2/|\mathbf{E}_0|^2$ can be obtained straightforwardly from the reflected fields given by Eqs. (36)–(38). It can be written as

$$R=R_1[\cos^2\beta\sin^2\varphi+R_2(\cos^2\beta\cos^2\varphi+\sin^2\beta)], \tag{39}$$

where

$$R_1=\left|\frac{\sqrt{\omega^2\cos^2\alpha-\omega_p^2}-\omega\cos\alpha}{\sqrt{\omega^2\cos^2\alpha-\omega_p^2}+\omega\cos\alpha}\right|^2 \tag{40}$$

and

$$R_2=\left|\frac{\cos\alpha\sqrt{\omega^2\cos^2\alpha-\omega_p^2}-\omega\sin^2\alpha}{\cos\alpha\sqrt{\omega^2\cos^2\alpha-\omega_p^2}+\omega\sin^2\alpha}\right|^2. \tag{41}$$

The first term in the *rhs* of Eq. (39) corresponds to $\beta=0$ ($\varphi=\pi/2$; *s*-wave, electric field perpendicular to the plane of incidence), while the second term corresponds to $\beta=\alpha$ ($\varphi=0$; *p*-wave, electric field in the plane of incidence). It is easy to see that there exists a cusp (shoulder) in the behaviour of the function $R(\omega)$, occurring at the transparency edge $\omega=\omega_p/\cos\alpha$, where the reflection coefficient exhibits a sudden enhancement on passing from the propagating regime to the damped one, as expected (total reflection). The condition for total reflection can also be written as $\sin\alpha=\sqrt{\varepsilon}$, where $R=1$ ($R_{1,2}=1$), as it is well known. For illustration, the reflection coefficient is shown in Fig. 2 for $\beta=\pi/6$ and various incidence angles. The reflection coefficient is vanishing at $\omega^2=\omega_p^2/(1-\tan^2\alpha)$ for $\alpha=\beta<\pi/4$ ($R_2=0,\varphi=0$).

5. Plasma slab

We consider an ideal plasma slab of thickness d, extending over the region $0<z<d$ and bounded by the vacuum. The displacement field $\mathbf{u}$ can be represented as $(\mathbf{v},u_3)[\theta(z)-\theta(z-d)]$, where $\mathbf{v}$ is the displacement component in the (x,y)-plane and u_3 is the displacement component along the z-direction. The approach presented above for a semi-infinite plasma can easily be extended to this case. The analogous of the equation of motion (4) exhibits now two polarization contributions, arising from the two surfaces. The dielectric response similar to Eq. (9) is given by

$$\begin{aligned}\mathbf{v}=&\frac{i e\mathbf{k}\omega_p^2}{m}\cdot\frac{(2\omega^2-\omega_p^2)\Phi_0-\omega_p^2\Phi_d e^{-kd}}{(\omega^2-\omega_p^2)[2\omega^2-\omega_p^2(1-e^{-kd})][2\omega^2-\omega_p^2(1+e^{-kd})]}e^{-kz}\\&+\frac{i e\mathbf{k}\omega_p^2}{m}\cdot\frac{(2\omega^2-\omega_p^2)\Phi_d-\omega_p^2\Phi_0 e^{-kd}}{(\omega^2-\omega_p^2)[2\omega^2-\omega_p^2(1-e^{-kd})][2\omega^2-\omega_p^2(1+e^{-kd})]}\\&\cdot e^{kz-kd}-\frac{i e\mathbf{k}}{m}\frac{\Phi}{\omega^2-\omega_p^2}\end{aligned} \tag{42}$$

and $iku_3=\partial v/\partial z$, where $\Phi_0=\Phi(z=0)$, $\Phi_d=\Phi(z=d)$, $0<z<d$. The electric field is given by $E_\perp=4\pi nev$ and $E_\parallel=4\pi neu_3$. One can see that, beside the bulk plasmon mode ω_p^2, there appears two surface modes given by $\omega_p^2(1\pm e^{-kd})/2$, as it is well known. For $d\to\infty$ Eq. (42) becomes the first Eq. (9) for the semi-infinite plasma. For $d\to 0$ we get the well-known plasma frequency $\sqrt{(2\pi n_s e^2/m)k}$ for a sheet with surface electron density $n_s=nd$.

The bulk contribution to the energy loss is the same as for the semi-infinite plasma. We compute the surface contribution to the electron energy loss for $kd\gg\omega_p d/v\gg 1$, i.e. for a fast electron moving with velocity v, which, however, spends enough time in the sample to excite plasmons. For normal incidence the surface contribution consists of two oscillatory terms

$$\begin{aligned}P_s=&-\frac{e^2\omega_p}{vt}(\sqrt{2}\sin\omega_p t/\sqrt{2}-\sin\omega_p t)\\&-\frac{e^2\omega_p}{d-vt}[\sqrt{2}\sin\omega_p(d/v-t)/\sqrt{2}-\sin\omega_p(d/v-t)]\end{aligned} \tag{43}$$

corresponding to the two surfaces, for $0<t<d/v$. The total energy loss during the passage through the slab is given by

$$\int_0^{d/v}dtP_s\simeq\int_0^\infty dtP_s=-\pi(\sqrt{2}-1)\frac{e^2\omega_p}{v}. \tag{44}$$

We use again the equation of motion (18) and the retarded potentials given by Eqs. (19) and (20) in order to get the refracted field (field inside the slab), reflected ($z<0$) and transmitted ($z>d$) fields. The polarization field is given by the same equations (23), where the z-integration is limited to the region $0<z<d$. The same holds for the equations of motion (25). We solve these equations by the same method used above. Within the slab we have two waves of the form $e^{\pm i\kappa' z}$, one being the refracted wave through the first surface ($z=0$), the other being the reflected wave on the second surface ($z=d$). The wavevector κ' is given by the same Eq. (29) and the transparency edge is given by the same condition $\omega\cos\alpha=\omega_p$ as for a semi-infinite plasma. We get

$$v_2=A_2\left[e^{i\kappa' z}-\frac{\kappa-\kappa'}{\kappa+\kappa'}e^{2i\kappa' d}\cdot e^{-i\kappa' z}\right], \tag{45}$$

where

$$A_2=\frac{2eE_{02}}{m\omega_p^2}\cdot\frac{\kappa(\kappa-\kappa')(\kappa+\kappa')^2}{K^2[(\kappa+\kappa')^2-(\kappa-\kappa')^2e^{2i\kappa' d}]} \tag{46}$$

and

$$v_1=A_1\left[e^{i\kappa' z}-\frac{\kappa-\kappa'}{\kappa+\kappa'}\cdot\frac{\kappa\kappa'-k^2}{\kappa\kappa'+k^2}e^{2i\kappa' d}\cdot e^{-i\kappa' z}\right], \tag{47}$$

where

$$A_1=\frac{2eE_{01}}{m\omega_p^2}\cdot\frac{\kappa'(\kappa-\kappa')(\kappa+\kappa')^2(\kappa\kappa'+k^2)}{(\kappa+\kappa')^2(\kappa\kappa'+k^2)^2-(\kappa-\kappa')^2(\kappa\kappa'-k^2)^2e^{2i\kappa' d}}, \tag{48}$$

the third component can be obtained from $\kappa'^2u_3=ik(\partial v_1/\partial z)$. One can check the transversality of these waves and can compute the dispersion relations for the eigenvalues (bulk and surface plasmon-polaritons) in the like manner as for the semi-infinite plasma.

The reflected field is given by

$$E_1 = E_{01}(1 - e^{2i\kappa' d}) \times \frac{(\kappa^2 - \kappa'^2)(\kappa^2\kappa'^2 - k^4)}{(\kappa + \kappa')^2(\kappa\kappa' + k^2)^2 - (\kappa - \kappa')^2(\kappa\kappa' - k^2)^2 e^{2i\kappa' d}} e^{-i\kappa z},$$

$$E_2 = E_{02}(1 - e^{2i\kappa' d}) \frac{\kappa^2 - \kappa'^2}{(\kappa + \kappa')^2 - (\kappa - \kappa')^2 e^{2i\kappa' d}} e^{-i\kappa z} \qquad (49)$$

and $E_3 = -E_{03}(E_1/E_{01})$.

From the above results one can check the continuity of the electric field and electric displacement as well as the angle of total polarization given by $\tan^2\alpha = 1 - \omega_p^2/\omega^2 = \varepsilon$. If we take formally $e^{2i\kappa' d} \to 0$ we recover all the fields for the semi-infinite plasma. Indeed, for the semi-infinite plasma all the integrations to $z \to \infty$ are taken by assuming a vanishing factor $e^{-\mu z}$, $\mu > 0$, and letting μ go to zero. If we preserve this factor for the slab, it gives rise to factors of the form $e^{2i\kappa' d} e^{-\mu d}$, which are vanishing for $d \to \infty$. The limit $d \to 0$ (plasma sheet) cannot be taken directly on the above results ($\omega_p \sim 1/\sqrt{d}$, $\kappa' \sim i\omega_p/c$), because of the discontinuities arising from the θ-function. The calculations for a plasma sheet with a finite (superficial) charge density n_s must be done separately. They are left, together with other related results, for a forthcoming publication. The limit $\kappa' d \ll 1$ ($\kappa d \ll 1$) can be taken directly on the formulae given here. It corresponds to wavelengths much longer than the thickness of the slab.

The reflection coefficient for the plasma slab $R = |\mathbf{E}_{refl}|^2/|\mathbf{E}_0|^2$, where the reflected field is given by Eqs. (49), has a different structure than the reflection coefficient for the semi-infinite plasma. It can be written as

$$R = \frac{\omega_p^4}{c^4}|1 - e^{2i\kappa' d}|^2[R_1\cos^2\beta\sin^2\varphi + R_2(\cos^2\beta\cos^2\varphi + \sin^2\beta)], \qquad (50)$$

where

$$R_1 = \frac{1}{|(\kappa + \kappa')^2 - (\kappa - \kappa')^2 e^{2i\kappa' d}|^2} \qquad (51)$$

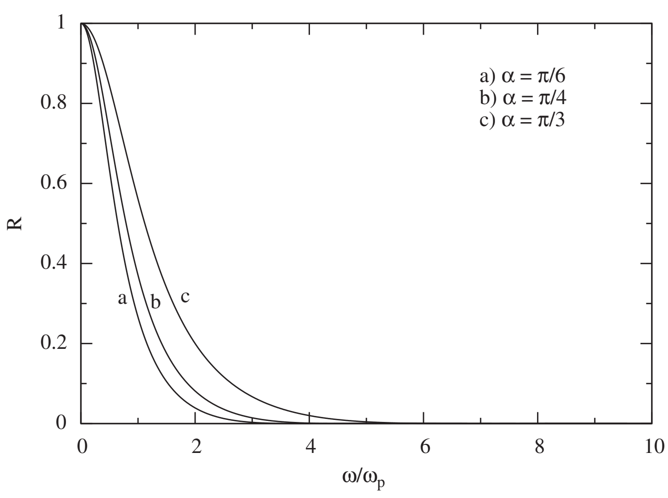

Fig. 3. Reflection coefficient for a slab of thickness d ($d\omega_p/c = 1$) for $\beta = 0$, $\varphi = \pi/2$ (s-wave) and a few incidence angles α. Its slope is continuous at the transparency edge ($\omega\cos\alpha = \omega_p$). The oscillations occurring in the transparency regime are too small to be visible in figure.

and

$$R_2 = \frac{|\kappa^2\kappa'^2 - k^4|^2}{|(\kappa + \kappa')^2(\kappa\kappa' + k^2)^2 - (\kappa - \kappa')^2(\kappa\kappa' - k^2)^2 e^{2i\kappa' d}|^2}. \qquad (52)$$

The reflection coefficient given by Eq. (50) is shown in Figs. 3 and 4 for $\beta = 0$, $\varphi = \pi/2$ (s-wave) and, respectively, $\alpha = \beta$, $\varphi = 0$ (p-wave) and $d\omega_p/c = 1$. The reflection coefficient exhibits characteristic oscillations arising from the exponential factor in Eqs. (50)–(52) and has an abrupt enhancement in the damping regime. In addition, R_2 is vanishing for $\omega^2 = \omega_p^2/(1 - \tan^2\alpha)$ ($\alpha < \pi/4$) and $R_2 = 1$ for $\omega = \omega_p$.

The transmitted field (region $z > d$) is given by

$$E_1 = E_{01}\frac{4K^2\kappa\kappa'(\kappa'^2 + k^2)e^{i(\kappa' - \kappa)d}}{(\kappa + \kappa')^2(\kappa\kappa' + k^2)^2 - (\kappa - \kappa')^2(\kappa\kappa' - k^2)^2 e^{2i\kappa' d}} e^{i\kappa z},$$

$$E_2 = E_{02}\frac{4\kappa'\kappa e^{i(\kappa' - \kappa)d}}{(\kappa + \kappa')^2 - (\kappa - \kappa')^2 e^{2i\kappa' d}} e^{i\kappa z} \qquad (53)$$

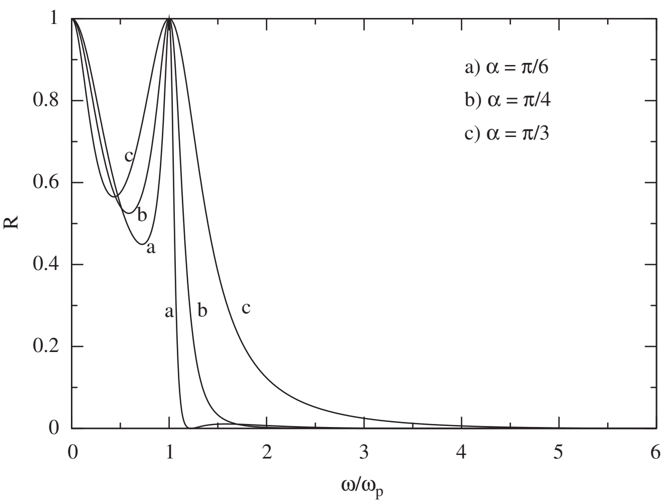

Fig. 4. Reflection coefficient for a slab of thickness d ($d\omega_p/c = 1$) for $\alpha = \beta$, $\varphi = 0$ (p-wave) and a few incidence angles α. It exhibits a local maximum ($R = 1$) for $\omega = \omega_p$ and small oscillations in the transparency region $\omega\cos\alpha > \omega_p$ (too small to be visible in figure). In addition, it is vanishing for $\omega^2 = \omega_p^2/(1 - \tan^2\alpha)$, $\alpha < \pi/4$, as one can see in figure for $\alpha = \pi/6$ (curve a).

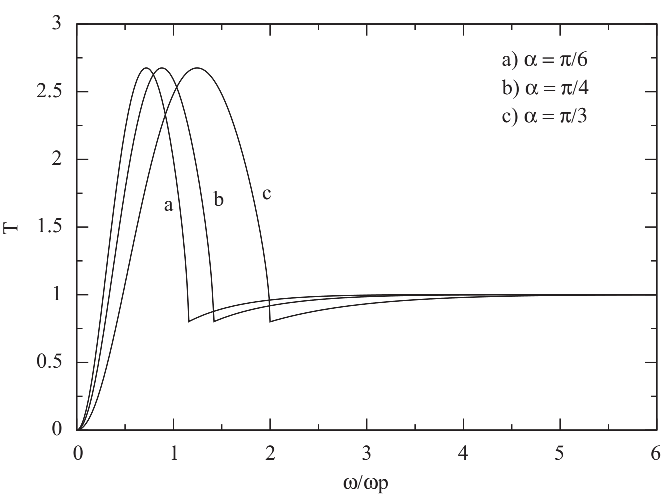

Fig. 5. Transmission coefficient for a slab of thickness d ($d\omega_p/c = 1$) for $\beta = 0$, $\varphi = \pi/2$ (s-wave) and a few incidence angles α. One can see the characteristic cusp at the transparency edge $\omega\cos\alpha = \omega_p$ and the peak occurring below this edge. The oscillations occurring in the transparency regime are too small to be visible in figure.

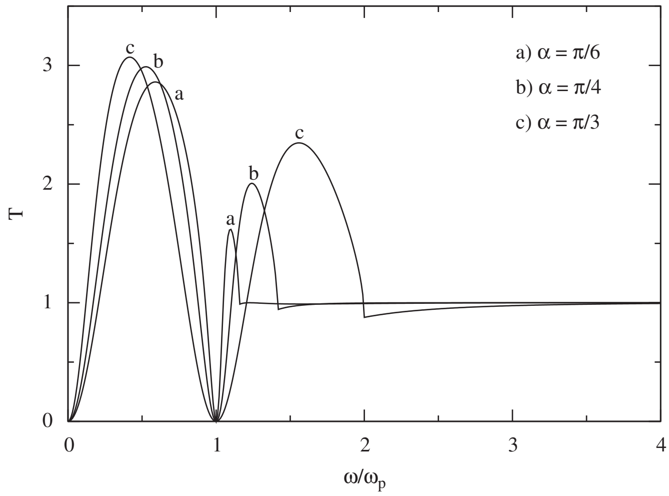

Fig. 6. Transmission coefficient for a slab of thickness d ($d\omega_p/c=1$) for a few incidence angles $\alpha=\beta$ and $\varphi=0$ (p-wave). One can see the two peaks occurring below the transparency edge $\omega\cos\alpha=\omega_p$ (the cusp in Figure) and the zero for $\omega=\omega_p$. The oscillations occurring in the transparency regime are too small to be visible in figure.

and $E_3=E_{03}(E_1/E_{01})$. One can check the continuity of the electric field and electric displacement at the surface $z=d$. In the limit $d\to\infty$ the transmitted field is vanishing. The transmission coefficient given by $T=|\mathbf{E}_{tr}|^2/|\mathbf{E}_0|^2$, where $\mathbf{E}_{tr}$ is given by Eqs. (53), can be written as

$$T=16\kappa^2|\kappa'|^2\left[R_1\cos^2\beta\sin^2\varphi+\frac{K^4|\kappa'^2+k^2|^2}{|\kappa^2\kappa'^2-k^4|^2}R_2(\cos^2\beta\cos^2\varphi+\sin^2\beta)\right], \tag{54}$$

where $R_{1,2}$ are given by Eqs. (51) and (52). This transmission coefficient is shown in Figs. 5 and 6 for $\beta=0$, $\varphi=\pi/2$ (s-wave) and, respectively, $\alpha=\beta$, $\varphi=0$ (p-wave) and $d\omega_p/c=1$. Beside the characteristic cusp occurring at the transparency edge ($\omega\cos\alpha=\omega_p$), the transmission coefficient exhibits an appreciable enhancement below this edge. For $\alpha=\beta$, $\varphi=0$ (p-wave) and $\omega=\omega_p$, the reflection coefficient attains the value unity and the transmission coefficient vanishes. The fields derived above can be viewed as generalized Fresnel's relations for a plasma slab.

6. Conclusions

The approach presented here is a quasi-classical one, valid for wavelengths much longer than the amplitude of the Fourier components of the displacement field $\mathbf{u}$. This is not a particularly restrictive condition for the classical dynamics of the electromagnetic field interacting with matter. When this condition is violated, as, for instance, for wavelengths much shorter than the mean separation distance between electrons, there appear both higher-order terms in the equations of motion and the coupling to the individual motion of the electrons. These couplings affect in general the dispersion relations and introduce a finite lifetime (damping) for the plasmon and polariton modes.

Making use of the equations of motion for the displacement field $\mathbf{u}$ and the radiation formulae for the electromagnetic potentials, we have computed herein the plasmon and polariton modes for an ideal semi-infinite electron plasma and an ideal plasma slab of finite thickness, as well as the dielectric response, the electron energy loss, the reflected and refracted waves and the reflection coefficient. For the semi-infinite plasma we have identified the bulk and surface plasmon–polariton modes and for the plasma slab we have computed also the transmitted wave and the transmission coefficient. It was shown that the stopping power due to the surface effects has a characteristic oscillatory behaviour in the transient regime near the surfaces. The field inside the plasma is either damped (evanescent) or propagating, as it is well known, and the reflection coefficient for the semi-infinite plasma exhibits a sudden enhancement on passing from the propagating to the damped regime, as expected. The transparency edge is given by $\omega\cos\alpha=\omega_p$, where α is the incidence angle, ω is the frequency of the incident wave and ω_p is the plasma frequency. Apart from characteristic oscillations, the reflection and transmission coefficients for the plasma slab exhibit an appreciable enhancement below the transparency edge.

Other effects related to the dynamics of a semi-infinite electron plasma, or, in general, various plasmas with rectangular geometries, can be computed similarly by using the method presented here. The method can also be applied to plasmas with other, more particular, geometries. The dissipation can be introduced (as for metals) and a model can be formulated for dielectrics, amenable to the method presented here. This will allow the treatment of more realistic cases as well as various interfaces, in particular plasmas (or metals) bounded by dielectrics. These investigations are left for forthcoming publications.

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