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LETTER TO THE EDITOR

Coulomb screening in the presence of two electromagnetic waves†

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Abstract. The effective static dielectric constant of an electron plasma in the presence of two electromagnetic waves is discussed. It is found that the potential distribution created by a point charge in such a system is anisotropic, dies off by power laws, and is enhanced when the difference of the two electromagnetic wave frequencies matches the plasma frequency.

There has been recent interest in the study of the interaction of intense laser fields with semiconductors (Ephstein 1970, 1971, Puchkov and Ephstein 1974, Bass and Granovskii 1971, Miranda 1976, Luzzi and Miranda 1978) and plasmas (Cohn *et al* 1972, Seely and Harris 1973, Amato and Miranda 1976). In the present work we look at another side of this problem, namely, the effect of electromagnetic waves on the screening of a static charge in an electron plasma.

As is well known, in the absence of an electromagnetic wave, the interaction of a long-wavelength longitudinal field (i.e., a sound wave or the field of an impurity) with conduction electrons in semiconductors is substantially lowered by the Coulomb screening factor. This is manifested, for example, in the acoustoelectric effect where the screening factor weakens the sound amplification by drifting electrons and one attempts to decrease this effect in various ways (Pomerantz 1964, Zilberman and Mishin 1970). In the presence of an electromagnetic wave, however, the space charge has high-frequency components at the wave frequency and its harmonics which, due to the nonlinearity of the system, contribute to the low-frequency components of the potential (Aliev and Silin 1965). Since the long-wavelength permittivity is nearly zero close to the plasma frequency ω_p , one might therefore expect a reduction of the weakening effect of screening on illuminating our sample with an electromagnetic wave of frequency ω near ω_p . This condition of $\omega = \omega_p$ imposes, in turn, severe restrictions from the experimental point of view. In the first place, $\omega = \omega_p$ is the threshold condition for the propagation of an electromagnetic wave in the plasma. Secondly, it is not easily achievable for the majority of semiconductor plasmas with the laser sources presently available. This drawback can, however, be avoided by using two transparent lasers of frequencies ω_1 and ω_2 ($\omega_1, \omega_2 \gg \omega_p$) such that $\omega_1 - \omega_2$ matches the plasma frequency, as will be shown below.

Let us consider a one-component plasma, infinite and homogeneous. Assume further that the condition $\omega_p \tau \gg 1$ is satisfied (i.e., one has a collisionless plasma). If the distance

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over which the amplitude of an electromagnetic wave changes is large in comparison with the size of the charge placed in the plasma, the initial Debye screening radius r_D , and the amplitude of the electron oscillations in the wave fields, we can use the dipole approximation. Hence, the Hamiltonian of our system can be written as ($\hbar = 1$)

$$H(t) = (1/2m) \sum_p [\mathbf{p} + (e/c) \mathbf{A}(t)]^2 c_p^+ C_p - e \sum_{\mathbf{p}, \mathbf{k}} \varphi(\mathbf{k}, t) c_{\mathbf{p}+\mathbf{k}}^+ c_{\mathbf{p}}, \quad (1)$$

where

$$\mathbf{A}(t) = (c/\omega_1) \mathbf{E}_1 \cos \omega_1 t + (c/\omega_2) \mathbf{E}_2 \cos \omega_2 t$$

describes the laser fields and the scalar potential describes the field of the static charge and the self-consistent field. The Fourier components of this scalar potential are given by the Poisson equation

$$K^2 \varphi(\mathbf{k}, t) = 4\pi \rho(\mathbf{k}) - 4\pi e \sum_p \langle c_{\mathbf{p}-\mathbf{k}}^+ c_{\mathbf{p}} \rangle_t, \quad (2)$$

where $\rho(\mathbf{k})$ is the Fourier component of the static charge, and $\langle \dots \rangle_t$ denotes averaging with the complete Hamiltonian.

Constructing the equation of motion for $\langle c_{\mathbf{p}-\mathbf{k}}^+ c_{\mathbf{p}} \rangle_t$ within the usual RPA (Pines 1961, Miranda 1973, 1975), solving it with the initial condition $\langle c_{\mathbf{p}-\mathbf{k}}^+ c_{\mathbf{p}} \rangle_{t=-\infty} = 0$, and substituting into (2), one gets

$$\varphi(\mathbf{k}, t) = \frac{4\pi \rho(\mathbf{k})}{k^2} - \frac{4\pi e^2}{k^2} \int_{-\infty}^t dt' \varphi(\mathbf{k}, t') \sum_p (f_{\mathbf{p}-\mathbf{k}} - f_{\mathbf{p}}) x$$

where $\mathbf{a}_i = e\mathbf{E}_i/m\omega_i^2$ is the electron oscillation amplitude in the field of wave i , $\epsilon_p = p^2/2m$, and f_p is the electron occupation number. If we now define

$$\begin{aligned} \tilde{\varphi}(\mathbf{k}, t) &= \varphi(\mathbf{k}, t) \exp(i\mathbf{k} \cdot \mathbf{a}_1 \sin \omega_1 t) \exp(i\mathbf{k} \cdot \mathbf{a}_2 \sin \omega_2 t) \\ \tilde{\rho}(\mathbf{k}, t) &= \rho(\mathbf{k}) \exp(i\mathbf{k} \cdot \mathbf{a}_1 \sin \omega_1 t) \exp(i\mathbf{k} \cdot \mathbf{a}_2 \sin \omega_2 t) \end{aligned} \quad (4)$$

and substitute them in (3), the relation between $\tilde{\varphi}$ and $\tilde{\rho}$ becomes the same as between φ and ρ in the absence of an electromagnetic wave. Hence, for the temporal Fourier components of $\tilde{\varphi}$ and $\tilde{\rho}$ we have (Pines 1961)

$$\tilde{\varphi}(\mathbf{k}, \omega) = 4\pi \tilde{\rho}(\mathbf{k}, \omega) / k^2 \epsilon(\mathbf{k}, \omega) \quad (5)$$

where $\epsilon(\mathbf{k}, \omega)$ is the usual dielectric constant in the RPA (Pines 1961). It then follows from (4) and (5) that

$$\begin{aligned} \varphi(\mathbf{k}, t) &= \exp(-i\mathbf{k} \cdot \mathbf{a}_1 \sin \omega_1 t) \exp(-i\mathbf{k} \cdot \mathbf{a}_2 \sin \omega_2 t) \\ &\times \int_{-\infty}^{+\infty} (d\omega/2\pi) [4\pi \rho(\mathbf{k}, \omega) / k^2 \epsilon(\mathbf{k}, \omega)] \exp(-i\omega t) \\ &\times \int_{-\infty}^{+\infty} dt' \exp(i\omega t') \exp(i\mathbf{k} \cdot \mathbf{a}_1 \sin \omega_1 t') \exp(i\mathbf{k} \cdot \mathbf{a}_2 \sin \omega_2 t'). \end{aligned} \quad (6)$$

Using the expansion

$$\exp(iz \sin \theta) = \sum_{n=-\infty}^{+\infty} J_n(z) \exp(in\theta)$$

in (6), and performing the indicated integrations one finally gets

$$\begin{aligned} \varphi(\mathbf{k}, t) = & \sum_{n, s, \mu, \nu = -\infty}^{+\infty} J_{n+\mu}(z_1) J_n(z_1) J_{s+\nu}(z_2) J_\nu(z_2) \\ & \times \{4\pi\rho(\mathbf{k}) \exp[i(\mu\omega_1 + \nu\omega_2)t]/k^2 \epsilon(\mathbf{k}, n\omega_1 + s\omega_2)\} \end{aligned} \quad (7)$$

where $z_i = \mathbf{k} \cdot \mathbf{a}_i$ ($i = 1, 2$). Equation (7) tells us that in the presence of high-frequency fields, the potential besides becoming anisotropic, has components at the field frequencies and their harmonics.

In the following, we consider only the static component $\varphi_0(\mathbf{r})$ of the potential (i.e., $\mu = \nu = 0$). We have from (7)

$$\varphi_0(\mathbf{r}) = 1/(2\pi)^2 \int d^3 k (4\pi\rho(\mathbf{k})/k^2 \epsilon_{\text{eff}}) \exp(-i\mathbf{k} \cdot \mathbf{r}) \quad (8)$$

where

$$1/\epsilon_{\text{eff}} = \sum_{n, s = -\infty}^{+\infty} J_n^2(\mathbf{k} \cdot \mathbf{a}_1) J_s^2(\mathbf{k} \cdot \mathbf{a}_2) / \epsilon(\mathbf{k}, n\omega_1 + s\omega_2) \quad (9)$$

i.e., the effect of two laser fields on the static potential can be taken into account by introducing an effective dielectric constant dependent on both the frequency and the polarisation of these fields. Notice that in the zero field limit ($a_i = 0$) only the $n = s = 0$ terms survive, so that ϵ_{eff} reduces to the usual static dielectric constant $\epsilon(\mathbf{k}, 0)$. The evaluation of (8) is, in general, quite difficult. However, assuming that the screened charge is a point static charge ($\rho(\mathbf{k}) = Ze$), and that the fields are such that $\omega_1 - \omega_2 = \eta\omega_p$, with η very close to unity, the situation is noticeably simplified. In this case, remembering that the high-frequency dielectric constant is $\epsilon \simeq 1 - \omega_p^2/\omega^2$, the dominant term in (9) is the resonant one, namely, $n = -s = 1$, with the result

$$\varphi_0(\mathbf{r}) \simeq [Ze/2\pi^2(\eta - 1)] \int d^3 k \exp(-i\mathbf{k} \cdot \mathbf{r}) \frac{J_1^2(\mathbf{k} \cdot \mathbf{a}_1) J_1^2(\mathbf{k} \cdot \mathbf{a}_2)}{k^2} \quad (10)$$

where we have approximated $\epsilon(\mathbf{k}, \omega_1 - \omega_2) = 1 - \omega_p^2/(\omega_1 - \omega_2)^2$, for $\omega_1 - \omega_2 = \eta\omega_p$, by $2(\eta - 1)$ assuming η is close to unity. It follows from (10) that in the presence of two laser fields with frequencies such that $\omega_1 - \omega_2$ is resonant with the plasma frequency, one has a screening breakdown resulting in an enhancement of the static potential.

As for the spatial variation of $\varphi_0(\mathbf{r})$, one may get a simple expression for not very strong fields ($ka_i \ll 1$) by expanding the Bessel functions in powers of the arguments, and assuming E_1 and E_2 circularly polarised. Under these circumstances (10) reduces to

$$\varphi_0(\mathbf{r}) \simeq \frac{3}{2} [Zea_1^2 a_2^2 / (\eta - 1) r^5] (\cos^2 \theta_1 - \frac{1}{2} \sin^2 \theta_1) (\cos^2 \theta_2 - \frac{1}{2} \sin^2 \theta_2) \quad (11)$$

where θ_1 and θ_2 are the angles of E_1 and E_2 with $\mathbf{r}$, respectively.

In conclusion, it should be emphasised that our discussion contains a number of simplified assumptions. Nevertheless, some essential conclusions can be drawn from it. Amongst them, the illumination of a semiconductor plasma with two laser fields ($\omega_1, \omega_2 \gg \omega_p$), such that $\omega_1 - \omega_2 = \omega_p$, can considerably change the screening properties of such plasma. For instance, even with an isotropic charge distribution, the potential turns out to be anisotropic, and die off by power laws which make the main contribution at large distances ($r \gg r_D$). Secondly, the polarisation of the electron cloud which screens a point charge is not of a dipole at large distance (Pipa 1970), but of a multipole nature.

Finally, when $\omega_1 - \omega_2 = \omega_p$ one has an enhancement of $\varphi_0(r)$ (screening breakdown). This enhancement of the potential could be experimentally observed, for example, by an increase in the electrical resistance of a semiconductor due to the scattering of conduction electrons by impurities. Furthermore, it may also play an important role in acoustoelectric devices as well as on the laser fusion problem—the enhancement of the electrical resistance implies an increase in the plasma heating.

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