Optical properties of crystals with spatial dispersion – Josephson plasma resonance in layered superconductors

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Abstract. – We derive the transmission coefficient, $T(\omega)$, for grazing incidence of crystals with spatial dispersion accounting for the excitation of multiple modes with different wave vectors $k$ for a given frequency $\omega$. The generalization of the Fresnel formulas contains the refraction indices of these modes as determined by the dielectric function $\epsilon(\omega, k)$. Near frequencies $\omega_e$, where the group velocity vanishes, $T(\omega)$ depends also on an additional parameter determined by the crystal microstructure. The transmission $T$ is significantly suppressed, if one of the excited modes is decaying into the crystal. We derive these features microscopically for the Josephson plasma resonance in layered superconductors.

Usually propagation of light in crystals is sensitive to average properties described by the dielectric function, but not to the specific atomic structure of the crystal, because the wave length of light is much larger than the interatomic distance. However, in crystals with spatial dispersion the atomic structure may affect optical properties at special extremal frequencies, $\omega_e$, where the group velocity vanishes, $v_g = \partial \omega(k) / \partial k$, of an eigenmode with dispersion $\omega(k)$ vanishes. Near $\omega_e$ the effective wave length, $\lambda_g = v_g / \omega$, becomes comparable with the interatomic distance, and then optical properties such as the reflectivity may feel the crystal microstructure.

Pekar [1] and Ginzburg [2], already realized that several eigenmodes with different wave vectors, $k$, at given $\omega$ may be excited by the incident light inside a crystal with spatial dispersion. Then the Maxwell boundary conditions (continuity of the components of the electric, $E$, and of the magnetic field, $H$, parallel to the surface), are not sufficient to find the relative amplitudes of these modes and hence the reflection coefficient. To overcome this problem within the macroscopic approach Pekar and Ginzburg introduced so called additional boundary conditions (ABC), which are supposed to be related somehow to the crystal microstructure near the surface. In this phenomenological approach the choice of these ABC is only heuristic and may be controversial, see ref. [3] and Comments to this paper. It is only a microscopic model which can determine the solutions inside the crystal unambiguously and demonstrate explicitly how the reflectivity depends on the crystal microstructure.

In this paper we present for the first time general macroscopic expressions for the transmission coefficient $T(\omega)$ in uniaxial crystals in the vicinity of $\omega_e$ to show how the Fresnel formula...
is modified by the presence of multiple solutions. Then we find a full microscopic description of optical properties for the Josephson plasma resonance (JPR) in layered superconductors, which supports the phenomenological picture and derives the ABC. We show explicitly that the maximum $T(\omega)$ reached at $\omega \approx \omega_c$ depends on the interlayer distance of the superconductor. Weak dissipation as well as a perfect crystal structure are necessary to observe this effect of spatial dispersion, as otherwise the influence of the microstructure is smeared out and the conventional Fresnel results are restored. For details of the calculation see [1].

We also point out that the stopping of light, i.e. $v_g = 0$, which in the case of gaseous media attracted considerable interest for the coherent optical information storage recently [3], here naturally appears in a solid at a finite wave vector due to spatial dispersion.

The JPR is an interlayer charge oscillation due to tunneling Cooper pairs and quasiparticles in layered superconductors. It is strongly underdamped because the quasiparticles responsible for dissipation are frozen out at low temperatures. The JPR is an appropriate phenomenon to study the effect of spatial dispersion on optical properties both theoretically and experimentally. Firstly, the quite simple Lawrence-Doniach model [4, 5] is sufficient to provide a complete microscopic description. Secondly, weak dissipation may be achieved at low temperatures if the JPR frequency is well below the superconducting gap. Thirdly, there is experimental evidence that the spatial dispersion of the Josephson plasmon is significant. Recently in the layered superconductor SmLa$_{1-x}$Sr$_x$CuO$_{4-\delta}$ two peaks at $\approx 7$ and $\approx 12$ cm$^{-1}$ were observed in optical properties [6]. These peaks can be naturally understood as the JPR of alternating intrinsic junctions with SmO or La$_1$Sr$_{0.5}$CuO$_{4-\delta}$, and the peak intensities indicate a strong inter-junction coupling [6, 7, 9–11], i.e. a significant spatial dispersion of the plasma modes. Tachiki et al. [6] mentioned the excitations of multiple modes and Marel et al. [9] suggested an effective dielectric function for alternating junctions, but in neither case the effects on the reflectivity presented below were discussed.

To formulate the general problem posed by spatial dispersion on a macroscopic level we consider a layered crystal with the dielectric functions $\epsilon_c(\omega)$ along the layers and $\epsilon_e(\omega, k_z)$ along the $z$-axis. In $\epsilon_c(\omega, k_z)$ we account for a collective mode (here JPR) with the dispersion $\omega_c(k_z)$, i.e. $\epsilon_c[\omega_c(k_z), k_z] = 0$. The incident polarized light of frequency $\omega$ has a magnetic field $\mathbf{B} = (0, B_y, 0)$ and a wave vector $\mathbf{k}_0 = (\omega \sin \theta/c, 0, \omega \cos \theta/c)$, where $\theta$ is the angle relative to the $z$-axis normal to the surface and to the layers. The bulk Maxwell equations

$$ck_zB_y = -\omega \epsilon_e(\omega, k_z)E_z, \quad ck_zB_y = \omega \epsilon_a(\omega)E_x, \quad k_xE_z - k_zE_x = -(\omega/c)B_y,$$

determine the wave vector $\mathbf{k} = (\omega \sin \theta/c, 0, k_z)$ of the eigenmodes at the given frequency $\omega$. Here $k_z(\omega)$ is a solution of the equation

$$k_z^2 = \epsilon_a(\omega)(\omega^2/c^2)[1 - \sin^2 \theta/\epsilon_c(\omega, k_z)].$$

When $\omega_c$ and $\epsilon_c$ are $k_z$-independent, eq. (2) has one solution for $k_z^2$ and the Maxwell boundary conditions lead to the conventional Fresnel formula

$$T = 1 - |1 - \kappa|^2/|1 + \kappa|^2 \quad \text{with} \quad \kappa = n/\epsilon_a(\omega)\cos \theta,$$

where $n = ck_z(\omega)/\omega$ is the refraction index. Then $T(\omega)$ is peaked near $\omega_c$.

In a crystal with spatial dispersion in $\epsilon_c(\omega, k_z)$, Eq. (2) has multiple solutions for $k_z^2(\omega)$, in the simplest case four (real or complex) ones, $\pm n_1(\omega), \pm n_2(\omega)$ [1, 2], which can interfere with each other [4]. For a semi-infinite crystal at $z > 0$ only two of them are physical. At low dissipation the energy flow (Poynting vector $\mathbf{S}$) of proper modes should be directed into the crystal in order to preserve causality. Thus their group velocity, $v_g(\omega) = \partial\omega(k_z)/\partial k_z \sim S_z$, should be positive [4]. In the case of normal (anomalous) dispersion this requires positive (negative) $\text{Re}(k_z)$. When dissipation is taken into account, this rule is equivalent to the condition that the eigenmodes decay inside the crystal, $\text{Im}(k_z) > 0$. Without dissipation one
Using eqs. (1)-(2) and (6), we derive relation for the amplitudes is determined in a microscopic model. Usually, Here \( \ell \) determine the amplitudes \( \gamma \) of the JPR, as it is confined between adjacent superconducting layers both in the bulk and on the surface. Generally this relation may differ near the surface. This difference is very small for the JPR, as it is confined between adjacent superconducting layers both in the bulk and on the surface, and the surface intrinsic junction is almost the same as that in the bulk. Therefore we restrict ourselves to the dielectric functions \( \epsilon_c(z, z') = \Theta(z)\Theta(z')\epsilon_c(z - z') \) with a sharp cutoff at the surface \( z = 0 \).

The ABC for our system, as proposed by Ginzburg,

\[
P_z(z) + \ell (\partial P_z / \partial z) = 0, \quad z \to 0,
\]
determine the amplitudes \( \gamma_{1,2} \) of the fields

\[
E_z(z) = \gamma_1 \exp(ik_{z1}z) + \gamma_2 \exp(ik_{z2}z), \quad P_z(z) = \gamma_1\chi_c(k_{z1}) \exp(ik_{z1}z) + \gamma_2\chi_c(k_{z2}) \exp(ik_{z2}z).
\]

Here \( \ell \) is a phenomenological parameter related to the crystal microstructure and is to be determined in a microscopic model. Usually, \( \chi_c(k_z) - \chi_c(0) \propto k_z^2 \) at \( k_z \to 0 \). In this limit the relation for the amplitudes is

\[
\gamma_1(1 + i\xi n_1) + \gamma_2(1 + i\xi n_2) = 0, \quad \xi = \omega\ell/c.
\]

Using eqs. (1)-(2) and (3), we derive

\[
\kappa = \frac{1}{\cos \theta} \frac{1 + \epsilon_a^{-1}n_1n_2}{n_1 + n_2 - i\xi\epsilon_c(1 + \epsilon_a^{-1}n_1n_2)}.
\]

In the limit \( |n_2| \gg |n_1| \gg \sqrt{|\epsilon_a|} \) the refraction index with smallest \( |n_p| \) determines \( \kappa \) and this condition defines the frequencies where the one-mode Fresnel description is correct. In the two-mode frequency interval \( T(\omega) \) is peaked when \( |n_1| = |n_2| \), and two situations are possible.

When both \( n_{1,2} \) are real without dissipation we have a frequency \( \omega_e \), where \( n_1 + n_2 = 0 \). At \( \omega = \omega_e \) only the term \( i\xi n_1n_2 \) with \( \xi \ll 1 \) remains in the denominator leading to \( T(\omega_e) = 0 \). As \( \omega \) increases above \( \omega_e \), the transmissivity increases and reaches its maximum,

\[
T_{\text{max}} = T(\omega_{e,\text{max}}) = 2/\{(1 + \epsilon_a^2\xi^2 \cos^2 \theta)^{1/2} + 1\},
\]

at \( \omega = \omega_{e,\text{max}} \) when \( n_1 + n_2 = (1 + \epsilon_a^{-1}n_1n_2)(\cos^{-2} \theta + \epsilon_a^2\xi^2)^{1/2} \) above \( \omega_e \). We see that \( T_{\text{max}} \) depends on \( \ell \) and is generally smaller than the Fresnel result \( T_{\text{max}} = 1 \) without dissipation. Note, that the opposite signs of the refraction indices \( n_{1,2} \) near \( \omega_e \) due to causality are essential for the dependence of \( T_{\text{max}} \) on the length \( \ell \). The vanishing of \( n_1 + n_2 \) at \( \omega_e \) in eq. (6) and its consequence on \( T_{\text{max}} \) have not been realized previously.

We anticipate an even stronger suppression of \( T_{\text{max}} \) due to spatial dispersion at frequencies \( \omega_i \) near \( \omega_{e_i} \), where \( n_1 \) is real, while \( n_2 = in_1 \) without dissipation. This occurs in superconductors when the dispersion of the collective mode is anomalous. \( T(\omega) \) is peaked at \( \omega_i \), but \( T(\omega_i) \ll T(\omega_{e,\text{max}}) \) (eq.(7)), as now \( n_1 + n_2 \neq 0 \) in the denominator of eq. (6). The microstructure \( \sim \xi \) is irrelevant at \( \omega_i \). This observation is confirmed below for the JPR.
The frequencies $\omega_c$ and $\omega_4$, the frequency interval of the two-mode regime near $\omega_c$ and $\omega_4$, as well as $n_{1,2}$ are determined by the dielectric function, but $T_{\text{max}}$ near $\omega_c$ and $\omega_4$ also depend on the ABC. These features and the ABC will now be derived microscopically for the JPR.

This requires the solution of the Maxwell equations and the equations for the phase of the superconducting order parameter, exactly accounting for the atomic (layered) structure of the crystal along the $z$-axis, supercurrents inside the 2D layers at $z = ms$ and for interlayer Josephson and quasiparticle currents between neighboring layers. We obtain

\[ c\partial_z B_y = i\epsilon_{d0}\sqrt{E_x - \frac{\omega_0^2}{\omega^2} \sum_{m=0} E_z s\delta(z - ms)} , \]  
\[ \partial_z E_x - ik_z E_z = i\frac{\omega}{c} B_y, E_{z,m,m+1} = \int_{ms}^{(m+1)s} E_z \frac{dz}{s} , \]  
\[ c\partial_z B_y = -\omega\epsilon_{d0} \left[ E_z - \sum_{m=0} A_m f_{m,m+1}(z) \right] , \]  
\[ \frac{\omega_0^2}{\omega_s^2} A_m = V_{m,m+1} = c s E_{z,m,m+1} + \mu_{m+1} - \mu_m . \]

Here $\mu_m$ is the chemical potential in the layer $m$, $V_{m,m+1}$ is the difference of the electrochemical potentials, $\omega_0^2 = 8\pi^2 c J_0 / \epsilon_{d0}$ are the bare JPR plasma frequencies determined by the Josephson critical current densities $J_0$ in junctions of type $l = 1, 2$, $\omega_0 = c / \lambda_{ab} \sqrt{\epsilon_{ab}}$ is the in-plane plasma frequency, $\epsilon_{d0}$ is the high frequency in-plane dielectric constant and $\lambda_{ab}$ is the in-plane London penetration length. The function $f$ is defined as $f_{m,m+1}(z) = 1$ at $ms < z < (m+1)s$ and zero elsewhere. To obtain eq. (11), for small amplitude oscillations we expressed the supercurrent density $J_{m,m+1} = J_l \sin \varphi_{m,m+1} \approx J_l \varphi_{m,m+1}$ via the phase difference $\varphi_{m,m+1} = 2iV_{m,m+1} / \hbar \omega$. Further, $\tilde{\omega}_c^2 = \omega^2 (1 - 4\pi\sigma_\text{cl} \omega / \omega_0^2 \epsilon_{d0})^{-1}$ takes into account the dissipation due to quasiparticle tunneling currents, $j_{m,m+1}^{(qp)} = \sigma_d V_{m,m+1} / \epsilon_s$, determined by the conductivities $\sigma_d$. We express the difference $\mu_m - \mu_{m+1}$ via the difference of the 2D charge densities, $\rho_m$, as $\mu_m = \rho_m / \rho_{m+1} = (4\pi\sigma_\text{cl} \omega / \omega_0^2 \epsilon_{d0}) (\mu_m / \mu_{m+1})$, where the parameter $\alpha = (\epsilon_{d0}/4\pi\epsilon s)(\partial \mu / \partial \rho)$ characterizes the inter-junction coupling, i.e. the JPR charge dispersion $\tilde{\omega}_c$. This parameter $\alpha$ is estimated as 0.4 for SmLa$_{1-x}$Sr$_x$CuO$_{4-\delta}$ [10]. Finally, the Poisson equation expresses $\epsilon_{d0}$ via $E_z$.

The solution of eqs. (11) - (12) between the layers $m$ and $m + 1$ is

\[ E_z(z) = c_m \exp(igz) + d_m \exp(-igz) + A_m , \]

and similar for $B_y$ and $E_x$, where $g = (\omega_0^2 \epsilon_{d0} / c^2 a)^{1/2}$ and $a^{-1} = 1 - \sin^2 \theta / \epsilon_{d0}$. The continuity of $B_y$ and $E_x$ across the layers leads to a set of finite-difference equations for $c_m, d_m$ and $A_m$.

First we consider a crystal with identical junctions, $\omega_0 = \omega_0 = c / \lambda_{ab} \sqrt{\epsilon_{d0}}$ and $\sigma_d = \sigma_c$. Omitting the terms of order $(gs)^2 \approx (s / \lambda_c)^2 \sim 10^{-10}$ we obtain in the bulk

\[ A_m(w - a - 2\alpha) + \alpha (A_{m+1} + A_{m-1}) = (1 - 2\alpha \beta) (c_m + d_m) , \]

for $A_m$ at $m \geq 1$ and similar equations for $c_m, d_m$. Here $\tilde{w} = \tilde{\omega}_c^2 / \omega_0^2$ and $\beta = s^2 / 2\lambda_{ab}^2 a$ ($\sim 10^{-4}$ in cuprates). Eliminating $c_m, d_m$, we find linear equations for $A_m$ alone, which give the dispersion of the eigenmodes. The difference between the equations for $A_0$ and $A_1$ is the microscopic boundary condition which allows us to find $A_m$ using bulk equations. The dispersion of the eigenmodes ($0 \leq q \leq 2\pi$),

\[ w(q) = \frac{\tilde{\omega}_c^2 q}{\omega_0^2} = 1 + 2\alpha (1 - \cos q) + \frac{(a - 1)\beta}{\beta + 1 - \cos q} , \]
Fig. 1 – Schematic picture of the dispersion $w = \omega^2/\omega_0^2$ for two alternating junctions ($\sigma_c = 0$, solid line). The lower band is analogous to the case of identical junctions and its dispersion is normal for $\theta = 0$ (dashed). Its mixing with a decaying electromagnetic wave (as shown by the dashed, vertical line at $\sin^2(q/2) = -2\beta$) results in two propagating modes with real $q$ near the lower band edge $w_c$, where the group velocity vanishes. The anomalous dispersion in the upper band gives rise to one propagating and one decaying mode and a special point $w_i$, where $\nu_1^2 = \nu_2^2$. is equivalent to eq. (2) with

\begin{equation}
\epsilon_c(\omega, q) = \epsilon_\infty \left[1 - \omega_c^2(q)/\omega^2 + 4\pi i \sigma_c/\omega \epsilon_\infty\right], \quad \epsilon_a(\omega) = \epsilon_\infty \left[1 - \omega_a^2/\omega^2\right],
\end{equation}

where $\omega_c^2(q) = \omega_\infty^2[1 + \alpha s^2 k_z^2]$ and $k_z^2 = 2(1 - \cos q)/s^2$. The dispersion of the plasma mode due to the charging effect, $\omega_p(q)$, is normal at low $q$ [6]. The last term in eq. (15) describes the anomalous (at small $q$) dispersion due to the inductive coupling of the in-plane currents [7]. When alone, it describes the decay of an electromagnetic wave inside the superconductor on the length $\lambda_{ab}$. As shown in fig. 1 for the lower band, their combination leads to the extremal frequency, $w_e = 1 + u$, with $u = [8(a - 1)/\beta\alpha]^{1/2}$, if $\alpha > (a - 1)/8$ and if the dissipation is weak, $4\pi \sigma_c/\omega \epsilon_\infty \ll u$. For any angle $\theta \neq 0$ the longitudinal plasma oscillations mix effectively with the electromagnetic wave at small $q \approx \sqrt{u/2\alpha}$. From eq. (16) we find $n_{1,2}^2 = (e/\omega s)^2[w - 1 \pm \sqrt{(w - 1)^2 - u^2}]/2\alpha$ and $|n_1 n_2| = \lambda_c^2 \epsilon_\infty u/2\alpha s^2 \approx \lambda_c^2/(s\lambda_{ab}) \gg 1$. Hence, the two-mode regime, when $|n_1| \sim |n_2|$, occurs in the frequency interval of width $\approx u \omega_0$ above $\omega_c$. Otherwise, at $|w - 1| \gg u$, the standard Fresnel formulas are correct.

Note that by modifying the JPR frequency $\omega_0$ and hence $\omega_c$, e.g. by an external magnetic field, the stopping of light at $\omega_c$ can be modified fastly, which might serve as the building block of a future magnetooptical device.

Next we find solutions for $A_m$, $c_m$, $d_m$. We obtain for $m \geq 1$

\begin{equation}
A_m = \gamma_1 A(q_1) \exp(iq_1 m) + \gamma_2 A(q_2) \exp(iq_2 m).
\end{equation}

The microscopic boundary condition found as described above is $A_{-1} = 0$, where $A_{-1}$ is the extension of the bulk solution, eq. (17), to $m = -1$. This corresponds simply to the fact that for the surface junction, $m = 0$, one neighboring junction, $m = -1$, is missing. We note that $A_m$ plays the role of the polarization, $P_z$, because it describes the response of the system to
the electrochemical potential, see eq. (12). With $A_{-1} = P_2(z = -s)$ we derive microscopically eq. (3) with $\ell = -s$. Then in eq. (8) we get $\xi \epsilon_a = s \lambda_c \sqrt{\epsilon_c / \lambda_{ab}}$ which may be of order unity in cuprates like Tl-2212 with $\lambda_c / \lambda_{ab} \sim 100$ and the JPR frequency $\sim 20 \text{ cm}^{-1}$.

Hence, for grazing incidence spatial dispersion ($\alpha \neq 0$) leads to (i) a shift of the peak position in $T(w)$ by $u \sim \sqrt{\beta}$, and (ii) a decrease of $T_{\text{max}}$ depending on the interlayer spacing $\sim \xi$, see fig. 2.

In a crystal with alternating Josephson junctions, $\delta = \omega_{01}^2 / \omega_{02}^2 < 1$, we introduce $A_1(n) = A_{2n}$ and $A_2(n) = A_{2n+1}$. They describe the polarization inside two different junctions in the unit cell $n$. For a plasmon decoupled from the electromagnetic field (at $\theta = 0$), the charge coupling of junctions leads to the dispersion of the two bands and a gap between them, see fig. 1. In the lower (upper) band the dispersion is normal (anomalous). Mixing with the decaying electromagnetic mode in the lower band leads to the existence of an extremal frequency $\omega_e$ for $\alpha^2 > (a - 1) \beta (1 - \delta) / 4$ and $4\pi \sigma q / \omega_{01} \epsilon_0 \ll u$. Hence, the situation in the lower band is similar to that in the case of identical junctions. In contrast to this, the mixing of the upper plasma band with anomalous dispersion and the decaying electromagnetic mode leads to an eigenmode with anomalous dispersion everywhere. Here $n_1$ is real, while $n_2$ is imaginary, and a critical frequency $\omega_i$ exists, see fig. 1.

When a junction of type 1 is at the boundary, the missing neighboring junction is of type 2 and the microscopic boundary condition is derived as $A_2(n = -1) = 0$. The solution of the equations for $A_1(n)$ and $A_2(n)$, i.e. the generalization of eqs. (14) and (17) for alternating junctions, are the eigenvectors $[A_1(q_p), A_2(q_p)]$, ($p = 1, 2$), which describe the microstructure of the local polarization within the unit cell for each mode $p$. Here we obtain

$$\kappa = \frac{1}{\cos \theta} \frac{[1 + \epsilon_a^{-1} n_1 n_2] [\hat{A}_2(q_1) n_2 - \hat{A}_2(q_2) n_1]}{A_2(q_1) n_2^2 - A_2(q_2) n_1^2},$$

where $\hat{A}_2(q) = A_2(q) \exp(-iq) \approx A_2(0)(1 - iq/2)$ at small $q = 2sk_z$. This leads in the lower band to eq. (7) for $\kappa$ and a similar behavior of $T(w)$ as in the case of identical junctions. In
the upper band the atomic scale $\xi$ is irrelevant. There the maximum value of $T(w)$ is reached at $w_{\text{max}} = w_i$ when $n_1 = in_2$ for $\sigma_{ql} = 0$,

$$T_{\text{max}} \approx \frac{2\alpha_{ab}^{3/2} \epsilon_{a0} \left(a - 1\right) L}{\lambda_c \left(se_{a0}\right)^{1/2} \cos \theta} \left(\frac{\alpha}{8\alpha a}\right)^{1/4}.$$  \hspace{1cm} (19)

Here $L = w_i(1 + \delta) - 2 - 8\alpha$, $w_i = (1 + \delta)(1 + 2\alpha)(1 + \sqrt{1 - p})/2\delta$ is the upper edge of the band at $a = 1$, $q = 0$ and $p = 4\delta(1 + 4\alpha)/(1 + \delta)^2(1 + 2\alpha)^2$. Thus $T_{\text{max}}$ depends on $\alpha$ explicitly and is smaller by the factor $(s/\lambda_{ab})^{1/2}$ than $T_{\text{max}}$ in the lower band or for the case with identical junctions, respectively. This difference vanishes in the Fresnel limit, when $4\pi\sigma_{ql} \gg \omega_0 u$.

In conclusion, multiple eigenmodes, propagating or decaying, are excited by incident light in crystals with spatial dispersion. At low dissipation this leads to a drop of the maximum transmission near resonance frequencies in comparison with the single-mode Fresnel result. In the case of two propagating modes near frequencies $\omega_c$, where the group velocity vanishes, the maximum transmission coefficient, $T_{\text{max}}$, depends on the crystal microstructure, which is in remarkable contrast to the general belief that atomic scales cannot affect optical properties due to the large wavelength of light. The drop of $T_{\text{max}}$ is larger when one excited mode is propagating, while the other is decaying. This behavior was demonstrated explicitly for the JPR, but it is general for any optically active excitations, e.g. optical phonons with anomalous dispersion in insulators. However, the condition of weak dissipation and a perfect crystal structure are crucial to observe deviations from the Fresnel regime.

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REFERENCES

[1] S.I. Pekar, Zh. Eksp. Teor. Fiz. 33,1022 (1957) [Soviet Phys. JETP 6,785 (1958)].
[2] V.M. Agranovich and V.L. Ginzburg, Spatial Dispersion in Crystal Optics and the Theory of Excitons, 2nd ed. (Nauka, Moscow, 1979) [Engl. transl. (Springer, Berlin, 1981)].
[3] K. Henneberger, Phys. Rev. Lett. 80,2889 (1998); ibid. 83, 1263.
[4] Ch. Helm, L.N. Bulaevskii, submitted to Phys. Rev. B.
[5] Chien Liu, et al., Nature 409 (2001) 490; D.F. Phillips, et al., Phys. Rev. Lett. 86,783 (2001).
[6] S.N. Artemenko et al., JETP Lett. 58,445 (1993), Physica C 253, 373 (1995); T. Koyama, et al., Phys. Rev. B 54, 16183 (1996), M. Tachiki, et al., ibid. 50, 7065 (1994).
[7] L.N. Bulaevskii, et al., Phys. Rev. B 50, 12831 (1994).
[8] H. Shibata, Phys. Rev. Lett. 86,2122 (2001); D. Dulic, et al., Phys. Rev. Lett. 86,4144 (2001); T. Kakeshita, et al., Phys. Rev. Lett. 86, 4140 (2001).
[9] D.v.d. Marel, et al., Phys. Rev. B 64, 024530 (2001).
[10] C. Helm, L.N. Bulaevskii, E.M. Chudnovsky, and M.P. Maley, cond-mat/0108449.
[11] Ch. Helm, et al., Physica C 362, 43 (2001); Ch. Preis, et al., SPIE Conf. Proc. 3480, San Diego, (1998), 236; D.A. Ryndyk, et al., J. Phys.: Cond. Mat. 14, (2002), 815. cond-mat/0108115.
[12] F. Forstmann and R.R. Gerhardts, Metal Optics near the Plasma Frequency, Springer tracts in modern physics vol. 109. Springer, Berlin, 1986.