COMPARATIVE ANALYSIS OF REFLECTION AND SURFACE POLARITON SPECTROSCOPIES AS METHODS FOR THE DETERMINATION OF INTRINSIC PLASMA FREQUENCIES IN SEMICONDUCTORS

E. F. Venger, A. V. Goncharenko, Yu. A. Pasechnik, and A. V. Mel'nychuk

The classical method for the determination of a plasma frequency, and thus the carrier concentration, which uses the minimum in the external plasma reflection spectrum [1-3] and a method which has recently become popular, based on surface polariton spectroscopy (and, in turn, the method of modified attenuated total internal reflection (ATIR) [4]) are convenient since they provide rapid analysis, are nondestructive to the sample, and do not require the preparation of contacts. At the same time it necessary to remember that these methods have some significant differences. Thus, surface polaritons of phonon type are permitted only in polar semiconductors, and the presence of doping often makes it necessary to consider a plasma-phonon interaction. Moreover, the localization depth for the surface polaritons may be significantly different from the thickness of the layer which contributes to the external reflection signal. The optimal conditions for the measurement of the necessary parameters are also different.

In this work, a comparative analysis of the precision of these methods for an idealized case, using a semi-infinite homogeneous medium as a model is carried out. Thus, uncertainties related to the unsuitability of this model for the real structure of a sample surface, and related to distortions caused by the response function of the monochromator, are not considered here.

The error in the determination of the plasma frequency from the experimentally measured reflection spectra or by ATIR may be represented as

$$\Delta \omega_p = \sum_i \left| \frac{\partial \sigma_i}{\partial \omega_p} \right|^{-1} \Delta \sigma_i,$$

(1)

where $\sigma_i$ are the linearly independent spectral parameters which are measured, and which depend on the plasma frequency. Foremost among these parameters, for the case of external reflection, are the frequency position and depth of the reflection minimum, while for ATIR spectra the half-widths are also important. The influence of inaccuracy in also the incident angle is considered separately below.

The model of the dielectric susceptibility of a semiconductor using the classical Helmholtz–Kettler form is described when a single dipole-active lattice vibration is active

$$\varepsilon(\omega) = \varepsilon_\infty \left( \frac{\omega_p^2 - \omega^2}{\omega_p^2 - \omega_0^2} - \frac{\omega_p^2}{\omega^2 + i\omega \gamma} \right),$$

(2)

where $\varepsilon_\infty$ is the high frequency dielectric susceptibility; $\omega_L$ and $\omega_T$ are the frequencies of the longitudinal and transverse optic phonons, respectively; $\gamma = 1/\tau$ is a phenomenological plasmon damping constant; $\tau$ is the relaxation time for a charge carrier pulse; and $\omega_p$ is the plasma frequency. Reflection of light at normal incidence gives

$$R(\omega) = \left| \frac{\sqrt{\varepsilon(\omega) - 1}}{\sqrt{\varepsilon(\omega) + 1}} \right|^2 \frac{z - \sqrt{2} [\varepsilon(\omega) + 1]}{z + \sqrt{2} [\varepsilon(\omega) - 1]},$$

(3)

where $\varepsilon(\omega) = \varepsilon_1(\omega) + i\varepsilon_2(\omega)$; and $z = \sqrt{\varepsilon_1(\omega) + \varepsilon_2(\omega)}$. In a case where the plasmon damping is very small and it may be neglected ($\gamma = 0$), then $R_{\min} = R(\omega_{\min}) = 0$ when $\varepsilon(\omega_{\min}) = 1$. From this we obtain
For strongly doped materials \((\omega_{\text{min}} > \omega_p)\) far from the region of residual beams or for nonpolar semiconductors it follows from (4) that

\[
\omega_p \approx \omega_{\text{min}} \sqrt{1 - \frac{1}{\varepsilon_\infty}}.
\]  

The presence of damping produces, as a rule \(R_{\text{min}} \neq 0\). For an estimate of the influence of damping on the reflection spectra, under the condition \(\omega_p \tau > 1\), we use the relationships

\[
\begin{align*}
\kappa^2_{\text{min}} &= 4R_{\text{min}}, \\
n_{\text{min}} &= 1 + 5R_{\text{min}},
\end{align*}
\]

where \(n_{\text{min}}\) and \(k_{\text{min}}\) are the values of the optical constants at the point of the minimum; \(n^2(\omega) - k^2(\omega) = \varepsilon_1(\omega)\); and \(2n(\omega)k(\omega) = \varepsilon_2(\omega)\). Then, from (2), we obtain

\[
\varepsilon_1(\omega) = \varepsilon_\infty \left( \frac{\omega^2 - \omega_{\text{min}}^2}{\omega^2 - \omega_p^2} - \frac{\omega_p^2}{\omega_{\text{min}}^2 + \tau^2} \right) = 1 + 6R_{\text{min}} + 25R_{\text{min}}^2.
\]  

A decomposition into a power series of the small quantity \(R_{\text{min}}\) under the condition \(\omega_p > \omega_\tau\) (strong doping), after straightforward but cumbersome calculations, yields

\[
\frac{\omega_{\text{min}}^2}{\omega_p^2} \approx 1 + \beta(\beta + 1) - R_{\text{min}} \left[ \frac{16}{\varepsilon_\infty} (1 - \beta) - \frac{6}{\varepsilon_\infty} \right] +
\]

\[
+ R_{\text{min}}^2 \left[ \frac{4}{\varepsilon_\infty} \left( 3 - \frac{8}{\varepsilon_\infty} \right) \left( 3 - \frac{16}{\varepsilon_\infty} \right) + \frac{25}{\varepsilon_\infty} (2\beta + 1) \right],
\]

where \(\beta = \frac{\omega^2 - \omega_p^2}{\omega_p^2} + \frac{1}{\varepsilon_\infty}\). Numerical estimates show that for the typical parameter values \(R_{\text{min}} \sim 0.1\) and \(\varepsilon_\infty \sim 10\) the derivative is

\[
\frac{\partial}{\partial R_{\text{min}}} \left( \frac{\omega_{\text{min}}}{\omega_p} \right) \sim \frac{1}{2\varepsilon_\infty} \left( 50R_{\text{min}} + 6 - \frac{16}{\varepsilon_\infty} \beta \right) \sim 0.5.
\]  

Thus, the shift of the frequency of the minimum in the plasma reflection as \(R_{\text{min}}\) changes is larger the larger \(R_{\text{min}}\) itself is. When \(R_{\text{min}} \sim 0.1\), if the precision of the reflection coefficient determination is \(\Delta R_{\text{min}} = 1\%), then this shift for semiconducting materials in the IR region may be several inverse centimeters, which thus exceeds the resolution of the apparatus. This indicates that for a precise determination of the plasma frequency from the external reflection spectra it is necessary to take damping into account, except perhaps in the case when \(R_{\text{min}} \leq 2\%\) (which corresponds, according to our estimates, to values of \(\gamma/\omega_p \leq 0.1\)). In the case of strong damping (when \(\gamma \geq \omega_p\)) the calculations show that the relative error \(\Delta \omega_p/\omega_p\) due solely to the imprecise determination of \(R_{\text{min}}\) may range from 1\% to 10\% or even higher. Moreover, knowing two parameters \((\omega_{\text{min}}\) and \(R_{\text{min}}\)), we may, in addition to the plasma frequency, determine \(\tau\) as well, which thus yields the charge carrier mobility, \(\mu = e\tau/m^*\).

The dispersion of the surface polaritons at the boundary between a semi-infinite medium and a vacuum is given by [5]

\[
k_x = \frac{\omega}{c} \sqrt{\frac{\varepsilon(\omega)}{\varepsilon(\omega) + 1}},
\]

where \(k_x\) is the component of the wave vector directed parallel to the surface. Experimentally, the surface polariton frequency (for relatively large gap values [6]) corresponds to the minimum in the ATIR spectrum. When the ATIR spectra are measured in the Otto geometry [5], \(k_x = (\omega/c)\text{nsin}\theta\), where \(n\) is the index of refraction of the ATIR hemicylinder, and \(\theta\) is the angle on incidence of the light on the hemicylinder. Expressing \(\varepsilon(\omega)\) in the form (2), we obtain the plasmon-phonon surface polaritons, and their dispersion relation for \(\gamma = 0\) may be written as