Generalization of the Landau–Placzek theory to media containing mobile charge carriers

V. A. Voĭtenko

St. Petersburg State Technical University, 195251 St. Petersburg, Russia
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A generalization of the Landau–Placzek theory, which is used to treat quasielastic light scattering spectra of media having a complicated and imperfect structure, is devised. Media containing mobile charge carriers are studied. A new theory of hydrodynamic fluctuations in such media is developed for this purpose. "Forced" solutions of a system of coupled inhomogeneous diffusion and heat-conduction equations corresponding to density and temperature fluctuations are obtained. Random heat and density fluxes serve as the fluctuation sources. The discovery that there are no correlations between them provides a substantial simplification. It allows the extinction coefficient for light scattering to be represented in the form of a simple rational function of the difference frequency. Numerical results of this representation are compared with experimental results for semiconductor crystals. An important parameter of matter, viz., the thermodiffusion coefficient D_T of the mobile charge carriers, is determined for the first time. The representation of the extinction coefficient found is also applicable to media containing mobile charge carriers which are not described by a kinetic equation. Experiments on quasielastic light scattering in superionic glasses are discussed as an example. © *1996 American Institute of Physics.* [S1063-7761(96)03004-1]

1. INTRODUCTION

Completely new solids containing mobile charge carriers, viz., compounds with an A_{15} structure, intermetallic actinides, high-T_c superconducting crystals, and superionic glasses, have been synthesized for research purposes. Sound undergoes strong dispersion in them, phonon anomalies associated with the charge carriers are observed, and the quasielastic light scattering spectra contain an intense wing adjacent to the laser line.[1,2] This wing has been thoroughly studied and attributed to electronic processes in semiconductors and plasmas.[3,4] Although electronic light scattering has not yet been observed in normal metals,[5] it has become a convenient tool for investigating intermetallic actinides with heavy fermions[6] and high-T_c superconducting crystals.[7] Quasielastic scattering has also been observed in superionic glasses[8] and has been used to investigate temperature waves.[9] The spectra of newly synthesized solids having a complicated and imperfect structure are usually described in terms of the macroscopic Landau–Placzek theory.[10,11] Thus, a need has arisen for a generalization of this theory to media containing mobile charge carriers with consideration of temperature fluctuations.

When mobile charge carriers are present in a medium, a self-consistent electric field E appears, which causes acoustic plasma oscillations.[12,13] They differ significantly from sound waves, i.e., isobaric density fluctuations, in molecular media (see, for example, Ref. 14). Therefore, the Landau–Placzek theory requires significant revision before it can be applied to solids.

Section 2 of the present work discusses the role of temperature fluctuations in quasielastic scattering. Particle-density fluctuations define a plasmon, i.e., shifted, scattering line. The central line is caused entirely by the temperature fluctuations. For this reason, the temperature (δT) and density (δn) fluctuations, rather than the entropy (δs) and pressure (δP) fluctuations traditionally used in the theory of molecular scattering,[11] were selected as the fundamental thermodynamic variables in this work. In Sec. 3 these fluctuations are studied under conditions with overlapping frequency ranges, i.e., for $\omega_p \tau \sim 1$, where ω_p is the characteristic frequency of the plasma oscillations and τ is their relaxation time. A new theory of hydrodynamic fluctuations taking place in a system of mobile charge carriers in condensed media under high-viscosity conditions is devised. An analytical expression for the scattering cross section, which describes manifestations of the interference of the "temperature waves" studied in Ref. 9 with plasma oscillations in spectra, is derived. In Sec. 4 the numerical results are compared with experimental results for semiconductor crystals and superionic glasses. It is shown that the theory devised faithfully describes the spectra. In addition, a rarely studied parameter of matter, viz., the thermodiffusion coefficient D_T of the mobile charge carriers, is determined for the first time. The calculation method developed in this work can be applied to the investigation of light scattering on hypersound with strong dispersion, which is not necessarily associated with free charge carriers.

2. LIGHT SCATTERING BY DENSITY AND TEMPERATURE FLUCTUATIONS

Light scattering is a process involving the emission of a secondary wave with a frequency ω_S, a wave vector $\mathbf{q}_S$, and a polarization vector $\mathbf{e}_S$ under the action of incident radiation with the parameters ω_I, $\mathbf{q}_I$, and $\mathbf{e}_I$. In the macroscopic theory it is convenient to express the extinction coefficient in

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terms of a correlator of the susceptibility fluctuations $\delta\chi(\mathbf{r},t)$ of the medium.[10] In the case of an isotropic medium, the corresponding expression has the form

$$\frac{\partial^2 h}{\partial\omega\partial\Omega}=\frac{1}{2\pi}\left(\frac{\omega_I}{c}\right)^4(\mathbf{e}_I\cdot\mathbf{e}_S)^2\int d\tau\int d^3r e^{i(\mathbf{k}\cdot\mathbf{r}-\omega\tau)}$$
$$\times\langle\delta\chi(\mathbf{R}+\mathbf{r},t+\tau)\,\delta\chi(\mathbf{R},t)\rangle. \qquad (1)$$

Here $\omega=\omega_I-\omega_S$, and $\mathbf{k}=\mathbf{q}_I-\mathbf{q}_S$. The general method for treating hydrodynamic fluctuations involves separating the statistically independent fluctuation parameters. For example, in the theory of molecular scattering the susceptibility fluctuation $\delta\chi(\mathbf{r},t)$ from (1) is expanded into a series in adiabatic pressure (sound) fluctuations, which relax in accordance with the viscosity of the medium, and isobaric entropy fluctuations, whose relaxation is governed by heat conduction. Such a treatment is applicable to Mandel'shtam–Brillouin scattering and the scattering of anisotropy fluctuations in molecular media.[10] However, in media containing free charge carriers, the set of fluctuation parameters must include the density, i.e., the concentration of quasiparticles δn. Since the thermodynamic parameter which is conjugate to the density is the temperature T, we find

$$\delta\chi=\left(\frac{\partial\chi}{\partial n}\right)_T\delta n+\left(\frac{\partial\chi}{\partial T}\right)_n\delta T. \qquad (2)$$

The macroscopic theory of scattering reduces to a calculation of the temporal correlation functions of the thermodynamic variables in terms of which the cross section is expressed after a expansion in the form (2) is substituted into (1). For molecular media without sound dispersion[15] this leads to the formula for the relative scattering intensity by isobaric entropy fluctuations and adiabatic sound fluctuations known as the Landau–Placzek ratio (see, for example, Refs. 10 and 11). Such a description of the various fine-structure components of the spectrum, which is based on their integrated intensities, is applicable only when the components are so narrow that their frequencies do not overlap (see, for example, Ref. 15). Nevertheless, an anomalously intense, broad structureless quasielastic scattering peak was observed in the intermetallic actinide UPt_3 in Ref. 16. It was attributed in the review in Ref. 6 to the enormous values of the Landau–Placzek ratio, $R=(C_P-C_V)/C_V=10^4-10^6$, where C_P and C_V are the specific heats at constant pressure and constant volume, respectively. Since Brillouin scattering lines were not observed at all in UPt_3, this "explanation" is essentially based on division by zero.

It has traditionally been assumed that scattering by the temperature fluctuations $\langle\delta T^2\rangle=T^2/C_V$ (Ref. 17), which can be described in terms of a total extinction coefficient of the form[10,11]

$$h=\frac{T^2\omega_I^4}{3c^4C_V}\left(\frac{\partial\chi}{\partial T}\right)_n^2, \qquad (3)$$

is difficult to observe because of its small amplitude. On the other hand, a quadratic temperature dependence of the integrated intensity corresponding to (3) was reliably recorded in the superionic glass $(AgI)_x(AgPO_3)_{1-x}$ for values of x from 0 to 0.55 (Ref. 8). A cubic dependence of the extinction coefficient on the temperature appears as a result of accidental degeneracy in some semiconductor crystals.[3,4,13] It seems likely that the scalar quasielastic scattering peak observed in Ref. 16 in the intermetallic actinide UPt_3 is also caused by temperature fluctuations. The spectral manifestations of temperature fluctuations just enumerated can be caused both by resonance enhancement and by the large specific heat of mobile charge carriers.[3,4]

The spectral composition of scattered light is determined by the kinetics of the susceptibility fluctuations $\delta\chi$ and therefore reflects the temporal correlation between the thermodynamic variables selected. Such a correlation also exists between independent parameters (see, for example, Ref. 14); therefore, when (2) is substituted into (1), all the correlation functions appearing and not just the mean squares must be considered. This gives

$$\frac{\partial^2 h}{\partial\omega\partial\Omega}=\frac{1}{2\pi}\left(\frac{e^2}{m^*c^2}\right)^2(\mathbf{e}_I\cdot\mathbf{e}_S)^2\{(\delta n^2)_{\mathbf{k}\omega}+r^2(\delta T^2)_{\mathbf{k}\omega}$$
$$+r[(\delta n\,\delta T)_{\mathbf{k}\omega}+(\delta n\,\delta T)_{-\mathbf{k},-\omega}]\}. \qquad (4)$$

Here m^* is the effective mass, the coefficient

$$r=\frac{(\partial\chi/\partial T)_n}{(\partial\chi/\partial n)_T} \qquad (5)$$

shows the degree of participation of the temperature fluctuations in the scattering, and the notations of the correlation functions are found in accordance with Ref. 10, for example,

$$(\delta n\,\delta T)_{\mathbf{k}\omega}=\int d\tau\int d^3r e^{i(\mathbf{kr}-\omega\tau)}$$
$$\times\langle\delta n(\mathbf{R}+\mathbf{r},t+\tau)\,\delta T(\mathbf{R},t)\rangle. \qquad (6)$$

To obtain the integrated scattering intensity, it is sufficient to switch from spectral to simultaneous correlation functions in (4). For this purpose, it should be assumed that the temporal arguments in (6) coincide, but since there is no simultaneous correlation of the independent parameters, this correlation function vanishes, and the terms remaining in (4) lead to the Einstein formula.[18] At the same time, the correlation function (6) can be significant for determining the spectral composition of scattered light, since its frequency dependence, which is calculated in the next section, has an antiresonance profile.

A similar picture is obtained for molecular scattering in the variables s and P. When there is strong sound dispersion in the medium, the Landau–Placzek theory by itself is insufficient. For example, in Sec. 6.8 in Ref. 19 the heat-conduction and Navier–Stokes equations were used to describe the scattering spectrum.

3. THEORY OF HYDRODYNAMIC DENSITY AND TEMPERATURE FLUCTUATIONS

In the media exhibiting scattering that were mentioned in the Introduction the charge carriers are characterized by low mobility. The macroscopic velocity is then equal to zero, and instead of the Navier–Stokes equation, the following diffusion equation must be used to determine the particle flux (see, for example, Sec. 59 in 20):

$$\delta\mathbf{j}=-D\nabla\delta n-\frac{D_T}{T}\nabla\delta T+\frac{\sigma}{e}\mathbf{E}+\mathbf{g}, \tag{7}$$

which should be supplemented by the continuity relation

$$\frac{\partial\delta n}{\partial t}+\operatorname{div}\,\delta\mathbf{j}=0. \tag{8}$$

Here D, D_T, and σ are the diffusion, thermodiffusion, and conductivity coefficients, and $\mathbf{g}$ is the random particle flux caused by thermal motion. The thermodiffusion current leads to the appearance of correlations between the density and the temperature, since it relates the diffusion equation (7) to the equation for the heat flux $\delta\mathbf{q}$ (see, for example, Sec. 59 in Ref. 20):

$$\delta\mathbf{q}-\mu\delta\mathbf{j}=\Pi\delta\mathbf{j}-\kappa\nabla\delta T+\xi. \tag{9}$$

Here μ is the chemical potential of the charge carriers,

$$\Pi=\frac{D_T}{D}\left(\frac{\partial\mu}{\partial n}\right)_T-T\left(\frac{\partial\mu}{\partial T}\right)_n \tag{10}$$

is the Peltier coefficient, κ is the thermal conductivity coefficient, and ξ is the random heat flux, which serves as the source of the temperature fluctuations. The continuity equation for the heat flux has the form

$$\frac{\partial\delta\varepsilon}{\partial t}+\operatorname{div}\,\delta\mathbf{q}=0, \tag{11}$$

where $\delta\varepsilon$ is the fluctuation of the volumetric energy density of the particles in the medium. The electric field $\mathbf{E}$ satisfies the Poisson equation

$$\varepsilon_0\operatorname{div}\mathbf{E}=4\pi e\delta n, \tag{12}$$

where ε_0 is the dielectric constant of the crystal lattice. For completeness, the system of hydrodynamic equations (7)–(9), (11), and (12) should be supplemented by an equation which relates $\delta\varepsilon$ to the fundamental thermodynamic variables δn and δT. For this purpose, in the expression for the derivative

$$\frac{\partial\delta\varepsilon}{\partial t}=C_V\frac{\partial\delta T}{\partial t}+\left(\frac{\partial\varepsilon}{\partial n}\right)_T\frac{\partial\delta n}{\partial t} \tag{13}$$

we employ the exact equation of state of a degenerate gas,[17] $P=2\varepsilon/3$, which gives

$$\left(\frac{\partial\varepsilon}{\partial n}\right)_T=\frac{3}{2}\left(\frac{\partial P}{\partial n}\right)_T=\frac{3}{2}n\left(\frac{\partial\mu}{\partial n}\right)_T. \tag{14}$$

As was shown in Ref. 15, the scattering spectrum is determined by the particular solution of the system of hydrodynamic equations associated with its inhomogeneous terms $\mathbf{g}(\mathbf{r},t)$ and $\xi(\mathbf{r},t)$. It is conveniently found by Fourier transforming Eqs. (7)–(9) and (11)–(13). Taking into account the frequency dispersion of the kinetic coefficients, for an unrestricted medium we obtain

$$\delta n=-i\frac{\mathbf{kg}(-i\omega+k^2\chi)-\dfrac{k^2D_T}{TC_V}\mathbf{k}\xi}{\left(-i\omega+k^2D+\dfrac{4\pi\sigma}{\varepsilon_0}\right)(-i\omega+k^2\chi)-\dfrac{i\omega k^2D_T^2}{TC_VD}\left(\dfrac{\partial\mu}{\partial n}\right)_T}, \tag{15}$$

$$\delta T=-\frac{i}{C_V}\frac{\mathbf{k}\xi\left(-i\omega+k^2D+\dfrac{4\pi\sigma}{\varepsilon_0}\right)-i\omega\dfrac{D_T}{D}\left(\dfrac{\partial\mu}{\partial n}\right)_T\mathbf{kg}}{\left(-i\omega+k^2D+\dfrac{4\pi\sigma}{\varepsilon_0}\right)(-i\omega+k^2\chi)-\dfrac{i\omega k^2D_T^2}{TC_VD}\left(\dfrac{\partial\mu}{\partial n}\right)_T}, \tag{16}$$

where $\chi=\kappa/C_V$ is the thermal diffusivity of the medium. The common denominator appearing in Eqs. (15) and (16) contains bands at imaginary frequencies, which correspond to temperature waves mixed with a plasmon (see also Ref. 9). Using (15) and (16), we can express each of the functions appearing in (4) in terms of quadratic functionals of $\mathbf{g}(\mathbf{r},t)$ and $\xi(\mathbf{r},t)$, whose mean values are determined by the correlation properties of these random fluxes. The correlator of the random currents g_i and g_l is found from the Nyquist relation (see, for example, Sec. 78 in Ref. 21)

$$(g_ig_l)_{k\omega}=\left(\frac{\partial n}{\partial\mu}\right)_T F(\omega)\delta_{il}\,\mathrm{Re}\,D(\omega), \tag{17}$$

where $F(\omega)=\hbar\omega[1-\exp(-\hbar\omega/T)]^{-1}$. The analogous correlator derived from the projections of the random heat fluxes ξ_i and ξ_l is also well known (see, for example, Sec. 88 in Ref. 21):

$$(\xi_i\xi_l)_{k\omega}=C_VTF(\omega)\delta_{il}\ \mathrm{Re}\ \chi(\omega). \tag{18}$$

The microscopic processes responsible for diffusion and heat conduction were considered in Refs. 3 and 4, as well as in Sec. 6.3 in Ref. 19, and it was shown that they are totally different in nature. A kinetic equation was used for this purpose in Refs. 3 and 4. For this reason, the random forces $\mathbf{g}$ and ξ corresponding to the diffusion fluxes $\delta\mathbf{j}$ and $\delta\mathbf{q}-(\mu-\Pi)\delta\mathbf{j}$ are uncorrelated, i.e.,

$$(\xi_ig_l)_{k\omega}=0. \tag{19}$$

Using (15) and (16) to write the spectral correlation functions of the concentration and the temperature, from (17)–(19) we find fairly elaborate explicit equations for them in terms of the dispersive kinetic coefficients. These equations are essentially reminiscent of the familiar Callen–Welton relation (see, for example, Ref. 17). Unfortunately, the character of the dispersion law for some of the kinetic coefficients appearing in (15) and (16) is unknown. Therefore, we shall not write out these cumbersome equations, noting only that the mixed correlation function (6) is proportional to $k^2D_T(\omega)$. To determine the spectral line shape from Eq. (4) we assume that the simplest form of the dispersion law compatible with the causality principle is valid:

$$\sigma(\omega)=\frac{\sigma_0}{1-i\omega\tau},\quad D(\omega)=\frac{D_0}{1-i\omega\tau}, \tag{20}$$

etc., where σ_0 is the static conductivity and D_0 is the corresponding diffusion coefficient. The conditions for the applicability of Eqs. (20) to systems described by a kinetic equation were considered in Refs. 3 and 4 and reduce to requiring that the collisions be elastic. Taking the dispersion law (20) for all the kinetic coefficients appearing in (15)–(18), for the cross section we obtain a rational function of the form

$$\frac{\partial^2 h}{\partial\omega\partial\Omega}=\frac{1}{2\pi}\left(\frac{e^2}{m^*c^2}\right)^2(\mathbf{e}_I\cdot\mathbf{e}_S)^2\left(\frac{\partial n}{\partial\mu}\right)_T\tau F(\omega)$$
$$\times\frac{NA-MB-(2N-M)(\omega\tau)^2+M(\omega\tau)^4}{[(\omega\tau)^4-(A+1)(\omega\tau)^2+B]^2+(\omega\tau)^2[A-2(\omega\tau)^2]^2}. \tag{21}$$

Here

$$A=(\omega_p\tau)^2+(\Gamma_1+\Gamma_2+\Gamma_C)\tau, \tag{22}$$

$$B=\Gamma_1\Gamma_2\tau^2+(\omega_p\tau)^2\Gamma_2\tau, \tag{23}$$

$$M=\Gamma_1\tau+2r_1\sqrt{\Gamma_1\Gamma_C}\tau+r_1^2(\Gamma_2+\Gamma_C)\tau, \tag{24}$$

$$N=\Gamma_1\Gamma_2\tau^2+r_1^2B, \tag{25}$$

where

$$\Gamma_C=k^2\frac{D_{T0}^2}{TD_0C_V}\left(\frac{\partial\mu}{\partial n}\right)_T, \tag{26}$$

$$\Gamma_1=k^2D_0, \quad \Gamma_2=k^2\chi_0, \tag{27}$$

$$r_1=r\sqrt{\left(\frac{\partial\mu}{\partial n}\right)_T\frac{T}{C_V}}. \tag{28}$$

Since Eq. (21) was obtained from general hydrodynamic diffusion and heat-conduction equations, the region where it is applicable is in no way related to the possibility of using the kinetic equation. The main result of the present work is the generalization of Eqs. (21)–(27) to media in which the kinetic equation is inapplicable: to crystals with a hopping conduction mechanism or with heavy fermions, high-T_c-superconducting crystals, and superionic glasses. Thus, the theory developed here is just as general as the Landau–Placzek theory. It adds the case of media containing free charge carriers to the latter.

4. NUMERICAL ANALYSIS OF THE FORM OF THE SPECTRUM AND COMPARISON WITH EXPERIMENT

Experiments on light scattering by mobile charge carriers in semiconductors using a Nd:YAG laser have a very long history. Despite the numerous publications of experimental scalar (polarized) scattering spectra,[22–26] the literature does not offer a single method for describing them over the entire frequency range. Equation (21) provides this possibility. To illustrate this we utilize the spectrum of Baĭramov and Toporov, which was borrowed from Ref. 22 and is represented by curve *1* in Fig. 1. Curve *2* shows the results of a calculation based on Eqs. (21)–(25) and the parameters listed in Table I. The values of Γ_1, Γ_2, ω_p, and $1/\tau$ were estimated preliminarily from experimental values of the concentration ($n=5.6\times10^{17}$ cm^{-3}) and the mobility ($b=2000$ cm^2/V·s) using the Einstein relation and the Wiedemann–Franz law (see the first row in Table I). The parameters r_1 and Γ_C are purely fitting parameters. The best agreement between theoretical curve *2* and experimental curve *1* is achieved for $r_1=0.85$ and $\Gamma_C=60$ cm^{-1}. The high value of r_1 detected is typical of many semiconductors. Owing to it, in particular, it was possible to observe light scattering by acoustic plasmons in both allowed[13] and forbidden[12] geometries.

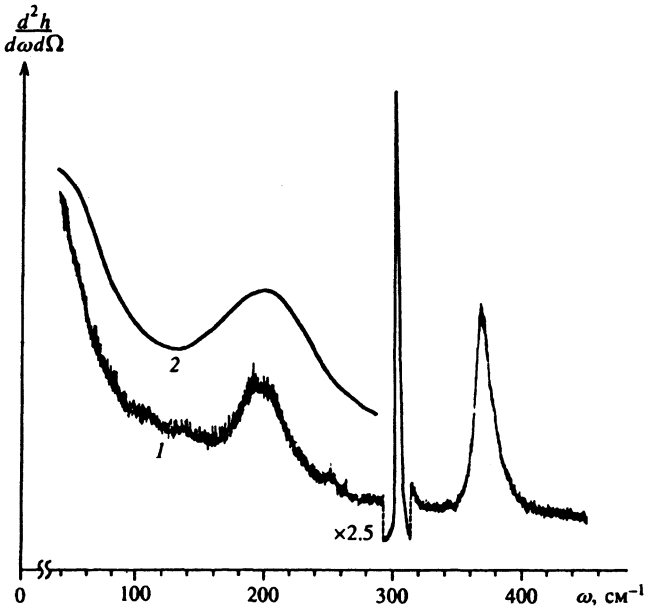

FIG. 1. Experimental quasielastic light scattering spectrum of a crystal of n-InP with $n=5.6\times10^{17}$ cm^{-3} at $T=300$ K (curve *1*) and its theoretical description according to Eq. (21) (curve *2*). The theoretical curve is shifted upward by an arbitrary amount for convenience in making the comparison. The spectra of light scattered on transverse and longitudinal optical phonons are also presented for calibration with indication of a scaling factor of 2.5.

The value obtained for the fitting parameter Γ_C makes it possible to determine the rarely studied thermodiffusion coefficient of mobile charge carriers in a substance. In the hydrodynamics of fluid mixtures (see, for example, Sec. 59 in Ref. 20) the thermodiffusion coefficient is understood to be the following linear function of $D_T(\omega)$:

$$D_T'(\omega)=\frac{D_T(\omega)}{n}-\frac{2TC_V}{n^2}\left(\frac{\partial n}{\partial\mu}\right)_T D(\omega). \tag{29}$$

TABLE I. Estimated and adjusted values of the parameters characterizing the spectra in Fig. 1.

	D_0, cm^2/s	χ_0, cm^2/s	D'_{T0}, cm^2/s	r_1	ω_p, cm^{-1}	$1/\tau$, cm^{-1}
Spectrum of n-InP (curve *1*) $n=5.6\times10^{17}$ cm^{-3} $b=2000$ cm^2/V·s $\mu=46$ meV	61[a]	47[b]	–	–	200	140
Calculated spectrum, (curve *2*)	59	81	113	0.85	200	115

[a]Evaluated from the Einstein relation.
[b]Evaluated from the Wiedemann–Franz law.

Although the sample of InP with $n=5.6\times10^{17}$ cm^{-3} used to record the spectrum (see curve *1* in Fig. 1) is in an intermediate stage of charge carrier degeneracy, we restrict ourselves here to the approximation for D_T' from (29), which is linear with respect to $T/\mu\ll1$. With this accuracy, from (26) we obtain

$$D_{T0}'=D_0\frac{T}{\mu}\pi\frac{2^{5/6}}{3^{4/3}}\sqrt{\frac{\Gamma_C}{\Gamma_1}}. \tag{30}$$

Substitution of the experimental values of Γ_1 and Γ_C from Table I gives the estimate $D_T'(\omega)\approx113$ cm^2/s.

The properties of the quasielastic scattering observed in the superionic glasses $(AgI)_x(AgPO_3)_{1-x}$ (Ref. 8) correspond to the properties considered in the present work. Its intensity increases linearly with the concentration x of the superionic component in accordance with (21) and is quadratically dependent on the temperature in accordance with (3), and its spectrum is a result of superposing two Lorentzian line shapes in accordance with the foregoing. A polarization dependence characteristic of quasielastic scattering by mobile changes in solids has been observed in other studies.[27,28] However, before the spectra could be identified completely, some difficulties arose in explaining the dependences of the width Γ_0 of the spectral line on the temperature and the concentration (i.e., on x). The classical theory of light scattering by mobile charges predicts only the possibility of a monotonic increase in Γ_0 with the temperature (see, for example, Fig. 5 in Ref. 13), although it allows both possibilities, i.e., an increase and a decrease in Γ_0 as a function of the mobility. No temperature dependence of Γ_0 was discovered in Ref. 8, in contradiction to the principles of the theory of scattering by mobile changes. Here it is assumed that $\Gamma_0=k^2D_0$. However, Eq. (21) allows another possibility, viz., $\Gamma_0=k^2\chi_0$. In this case the constancy of Γ_0 means only that the gas-dynamic equation $\chi_0=D_0$ is not applicable to media which have a superionic mechanism of conduction. It is possible, for example, that the strong temperature dependence of the specific heat C_V cancels the analogous dependence of the superionic thermal conductivity κ, so that, as a result, their ratio $\chi=\kappa/C_V$ does not depend on the temperature. Thus, the doubts expressed in Ref. 8 regarding the superionic nature of the quasielastic scattering spectrum are completely allayed.

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