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Optical properties of nonequilibrium low-dimensional systems

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The optical properties of low-dimensional carrier systems (‘‘quantum wire’’ type) driven away from equilibrium are studied. The frequency and wave-vector-dependent dielectric function of a quasi-one-dimensional electron system under the action of an exciting external pumping source is derived. The optical responses of the system are obtained in terms of its nonequilibrium thermodynamic state, the latter characterized resorting to a nonequilibrium statistical ensemble formalism.

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I. INTRODUCTION

The question of transport and other physical properties in spatially constrained systems is having a strong and healthy development. This applies to two kinds of areas of research, such as the physics of semiconductors and biological systems. The case of semiconductor systems with low dimensionality is presently the object of extensive theoretical and experimental study (e.g., see [1]); the case of biosystems involves among others, biopolymers, in particular, peptides and proteins, composed of long chains of amino acid units, and DNA molecules [2]. Both, semiconductors in electronic devices and living biological systems, work under nonequilibrium conditions (usually far from equilibrium), and then involve situations where strong dissipative effects are unfolding in the medium. Thus this question belongs to the realm of the thermodynamics of irreversible processes, which is a field theory at a macroscopic level dealing with states and processes in systems lying beyond equilibrium, either by a large amount or close to equilibrium corresponding to the nonlinear and linear, respectively, divisions of this area of research. Hence, nonequilibrium thermodynamics deals basically with transport phenomena involving changes in time and space of macroscopic observables and their fluxes in continuum media, as well as with the important case of steady states.

We consider here the case of a quasi-one-dimensional system (Q1DS) as a quantum wire semiconductor sample or a biopolymer containing charge carriers (we take the case of a *p*-doped material—where the carriers are holes—as it is, for example, the case in proteins [3]). We concentrate our analysis on the optical properties of these carriers, to evidence their electronic excitations (single particles and plasma waves) through the analysis of Raman scattering experiments under varying excitation conditions.

II. IRREVERSIBLE THERMODYNAMICS OF THE Q1DS

Let us consider a quasi-one-dimensional cylindrical sample, where mobile carriers are present—and we consider in particular the case of holes in *p*-doped materials—which are traveling in the ionic background. Let R be the radius of the wire and L its length; moreover, these carriers are excited by the action of an external source pumping energy on them. This exciting pump may be a source of electromagnetic radiation e.g., provided by a laser, in the case of the semiconductor (e.g., see [4]), or nonelectromagnetic radiation in the case of biopolymers (e.g., see [5]). Let n_h be the linear density of holes (number of holes per unit length along the axis of the wire), and in this cylindrical constrained geometry the hole energy levels are given by

$$\epsilon_{knl}=\frac{\hbar^2}{2m_{\mathrm{h}}^*}(\kappa_{nl}^2+k^2), \tag{1}$$

where m_{h}^* is the hole effective mass, k is the crystalline wave number for free movement along the z direction of the axis of the cylinder, and $\kappa_{nl}=\beta_{nl}/R$ where β_{nl} are the zeros of the Bessel function of order $n, J_n(\beta_{nl}r/R)$, with $n=0,1,2,\ldots$ and $l=1,2,\ldots$, and r is the radial coordinate (the problem is axially symmetric). The wave functions are in cylindrical coordinates

$$\Psi_{nlk}(r,\theta,z)=\mathcal{A}J_n(\beta_{nl}r/R)\exp(in\theta)\exp(ikz), \tag{2}$$

where

$$\mathcal{A}^{-1}=\sqrt{L}(\pi R^2)^{1/2}J_{n+1}(\beta_{nl}). \tag{3}$$

We recall that the knowledge of the frequency and wave-number-dependent dielectric function $\varepsilon(Q,\omega)$ provides information on all the optical properties of the system; here, we are interested in the particular one consisting of the Raman scattering spectrum when considering different conditions of excitation, that is, changing values of the intensity of the pumping source.

As already noticed, for nonequilibrium systems we are required to analyze their macroscopic properties using a thermomechanical statistical formalism, and we resort to the use of the nonequilibrium statistical operator method (NESOM) and Zubarev's approach is used [6,7]. According to NESOM, the first step—in this nonequilibrium statistical ensemble formalism—is the choice of a set of basic dynamical vari-

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ables (mechanical observables) that can provide a macroscopic description of the system appropriate for the experimental situation to be considered. The average values of these dynamical variables over the nonequilibrium ensemble provides the basic set of macrovariables that define the nonequilibrium thermodynamic space of states [8,9]. We consider a model where only the lowest subband $n=0$, $l=1$ [cf. Eq. (1)] is occupied, which then restricts the values of the doping concentration n_h we can use.

For the basic dynamical variables we take

$$\{\hat{H}_0\hat{N}_h;\{\hat{n}_{kQ}\}\}, \tag{4}$$

consisting of the Hamiltonian $\hat{H}_0$ and the operator $\hat{N}_h$ for the number of the free holes

$$\hat{H}_0=\sum_k \epsilon_k h^\dagger_{-k}h_{-k}, \quad \hat{N}_h=\sum_k h^\dagger_{-k}h_{-k}, \tag{5}$$

and the Dirac-Landau-Wigner single-particle dynamical operators for holes

$$\hat{n}_{kQ}=h_{-k-(1/2)Q}h^\dagger_{-k-(1/2)Q}, \tag{6}$$

where $Q\neq 0$, which is associated with the description of inhomogeneous quantities, while the diagonal terms (populations) are present in $\hat{H}_0$ and $\hat{N}_h$ of Eq. (5). Moreover, h ($h^\dagger$) are the usual annihilation (creation) operators in the quantum state labeled by the subindex k [we recall that we are only considering the subband $n=0$, $l=1$, indexes which we have omitted in Eqs. (4)–(6)].

Two points need to be stressed: the first is that in condensed matter physics, however the strong Coulomb interaction is present, a single-particle description in the random-phase approximation [11] can be used and it is extremely successful; the second is that the quantity of interest for us here is, as noticed, $\varepsilon(Q,\omega)$, which can be derived from the knowledge of the polarization charge (in units of the electron charge). In fact, from electrodynamic theory [10]

$$\frac{1}{\varepsilon(Q,\omega)}-1=\frac{n(Q,\omega)}{r_0}, \tag{7}$$

where $n(Q,\omega)$ is the Fourier transform of the charge density generated by a probe charge oscillating with frequency ω, propagating with wave number Q, and having amplitude r_0, which we calculate in the random-phase approximation.

The basic thermodynamic variables are then the averages over the nonequilibrium ensemble of the quantities of Eq. (4) at time t, which we designate as

$$\{E_0(t);n_h(t);\{n_{kQ}(t)\}\}, \tag{8}$$

which is the free holes energy and density and Dirac-Landau-Wigner single-particle density matrix in nonequilibrium conditions, namely,

$$E_0(t)=\mathrm{Tr}\{\hat{H}_0\varrho_\varepsilon(t)\}, \tag{9}$$

$$n_h(t)=\frac{1}{L}\mathrm{Tr}\{\hat{N}_h\varrho_\varepsilon(t)\}, \tag{10}$$

$$n_{kQ}(t)=\mathrm{Tr}\{\hat{n}_{kQ}\varrho_\varepsilon(t)\}, \tag{11}$$

where ϱ_ε is the nonequilibrium statistical operator. The latter is given in MaxEnt-NESOM [6,7,12,13] by the expression

$$\varrho_\varepsilon(t)=\exp\left\{-\hat{S}(t,0)+\int_{-\infty}^{t}dt'\,e^{\varepsilon(t'-t)}\frac{d}{dt'}\hat{S}(t',t'-t)\right\}, \tag{12}$$

where in this case, for the basic set of variables of Eq. (4),

$$\hat{S}(t,0)=\phi(t)+\beta(t)[\hat{H}_0-\mu(t)\hat{N}_h]+\sum_{kQ}F_{kQ}(t)\hat{n}_{kQ} \tag{13}$$

is the so-called informational entropy operator [14] and

$$\hat{S}(t',t'-t)=\exp\left\{-\frac{1}{i\hbar}(t'-t)\hat{H}\right\}\hat{S}(t',0)\exp\left\{\frac{1}{i\hbar}(t'-t)\hat{H}\right\}. \tag{14}$$

The quantity $\phi(t)$ (playing in the theory the role of the logarithm of a nonequilibrium partition function and being a nonequilibrium Massieu-Planck thermodynamic potential) together with $\beta(t)$, $-\beta(t)\mu(t)$, and $F_{kQ}(t)$ are the corresponding Lagrange multipliers that the variational derivation in MaxEnt-NESOM introduces.

The relevance of the choice of the set $\{n_{kQ}(t)\}$ resides in that we are looking for an expression for $\varepsilon(Q,\omega)$ of Eq. (7), where $n(\vec{Q},\omega)$ is the Fourier transform in time of

$$n(Q,t)=\sum_k n_{kQ}(t). \tag{15}$$

Leaving $E_0(t)$ for a later analysis, taking into account that $n_h(t)$ is the constant doping concentration n_h, using the nonlinear quantum kinetic theory that MaxEnt-NESOM provides [6,7,15] in the Markovian approximation [16], and considering that under the action of a constant-in-time source pumping energy in a uniform way, the system achieves a homogeneous steady state (after a transient regime has elapsed). The equations of evolution for Dirac-Landau-Wigner density matrices resulting from the perturbation of a Q- and ω-dependent probe charge are

$$\begin{aligned} i\hbar\frac{d}{dt}n_{kQ}(t)={}&r_0e^{i\omega t}(f^*_{k+(1/2)Q}-f^*_{k-(1/2)Q})\mathcal{V}(Q)\\ &+\Delta E_{kQ}n_{kQ}(t)-\mathcal{V}(Q)(f^*_{k+(1/2)Q}-f^*_{k-(1/2)Q})\\ &\times n(Q,t)+i\hbar[A_{kQ}+B_{kQ}]n_{kQ}(t)\\ &+i\hbar(f^*_{k+Q}-f^*_k)\sum_{k'}C_{k'Q}n_{k'Q}(t)\\ &+i\hbar\sum_{k'}D_{kk'Q}n_{k+k'Q}(t), \end{aligned} \tag{16}$$

where Q together with k and k' run over the Brillouin zone of the periodic array of molecules being considered. In this Eq. (16)

$$\Delta E_{kQ}=\epsilon_{k+(1/2)Q}-\epsilon_{k-(1/2)Q}, \tag{17}$$

and

$$\mathcal{V}(Q)=2e^2K_0(QR)/\epsilon_0 L \tag{18}$$

is the matrix element of Coulomb interaction for this Q1DS, with K_0 being the Bessel function of order zero and ϵ_0 is the background static dielectric constant. Moreover, the terms with *B*, *C*, and *D* involve contributions arising out of carrier-phonon interaction, and *A* the coupling with the external source, which we discuss below. The quantity

$$f_k^*=\mathrm{Tr}\{h_{-k}^\dagger h_{-k}\bar{\varrho}_{ss}\} \tag{19}$$

plays the role of a population in state $|k\rangle$ of the holes in the uniform steady state characterized by the statistical operator

$$\bar{\varrho}_{ss}=\exp\{-\phi^*-\beta^*\hat{H}_0+\beta^*\mu^*\hat{N}\}. \tag{20}$$

In this Eq. (20) we recall ϕ^* is the Lagrange multiplier which in NESOM ensures the normalization of the statistical operator, $\beta^*\equiv 1/k_BT_c^*$ is the one associated with the energy that has been written in the usual form introducing the carrier quasitemperature T_c^* and the quasichemical potential μ^* [17,18], which are time independent. Moreover, under constant uniform excitation, the Lagrange multipliers F_{kQ} are null in the steady state.

Furthermore, of the four coefficients *A*, *B*, *C*, and *D*, at intermediate to high excitations levels produced by the external pumping source, those associated with carrier-phonon interaction, namely, *B*, *C*, and *D*, are much smaller than the one due to *A* (responsible for the pumping effects) and are disregarded (and then we omit writing their cumbersome expressions). Hence, we take only the term with A_{kQ} which can be written in the form

$$A_{kQ}=g_{kQ}I_0, \tag{21}$$

where I_0 is the intensity of the source and g_{kQ} a coupling coefficient indicating the efficiency with which the pumped energy is absorbed. This coefficient depends on the type of external excitation that is used, and in what follows, without loss of generality, we take it to be independent of *k* and it is left as an open parameter in the theory. Finally, after some calculation we obtain that

$$f_k^*\simeq[1+\exp\{\beta^*(\epsilon_k-\mu^*)\}]^{-1}. \tag{22}$$

Furthermore, this distribution, in the usual experimental conditions, can be approximated by

$$f_k^*\simeq n_h\left(\frac{2\pi\hbar^2}{m_h^*k_BT_c^*}\right)^{1/2}\exp\left(-\frac{\epsilon_k}{k_BT_c^*}\right). \tag{23}$$

This follows from the condition that the exponential in Eq. (22) is much larger than 1, and for the given fixed number (density n_h) of holes, the relation [cf. Eq. (10)]

$$n_h=\frac{1}{L}\sum_k f_k^* \tag{24}$$

determines the coefficient in front of the exponential in Eq. (23); it is worth noticing that the population as given by Eq. (23) has an expression resembling a Maxwell-Boltzmann-like distribution for particles with density n_h and temperature T_c^*, and, we recall, *L* is the length of the sample. The quasitemperature is determined by the relation

$$\frac{1}{L}E_0=\frac{1}{L}\sum_k \epsilon_k f_k^*=\frac{1}{2}n_\mathrm{h}k_BT_c^*, \tag{25}$$

which follows after Eq. (23) is used, and where E_0 is the energy of the holes [cf. Eq. (9)] in the steady state attained when under the action of the constant external pump.

Hence, after the terms containing *B*, *C*, and *D* in Eq. (16) are neglected, we are left, after Fourier transformation in time is performed, with the integral equation

$$\begin{aligned}-\hbar\omega n_{kQ}(\omega)=&-r_0(f_{k+(1/2)Q}^*-f_{k-(1/2)Q}^*)\mathcal{V}(Q)\\&+\Delta E_{kQ}n_{kQ}(\omega)-\mathcal{V}(Q)\\&\times(f_{k+(1/2)Q}^*-f_{k-(1/2)Q}^*)n(Q,\omega)\\&+i\hbar A_Q n_{kQ}(\omega),\end{aligned} \tag{26}$$

where A_Q is the quantity of Eq. (21) taken as *k* independent and $n(Q,\omega)$ is given by

$$n(Q,\omega)=\sum_k n_{kQ}(\omega), \tag{27}$$

the quantity of Eq. (15) after taking Fourier transform in the time coordinate.

Equation (26) can be rewritten as

$$n_{kQ}(\omega)=r_0\mathcal{V}(Q)\mathcal{F}_{kQ}(\omega)+\mathcal{V}(Q)\mathcal{F}_{kQ}(\omega)n(Q,\omega), \tag{28}$$

where we have introduced

$$\mathcal{F}_{kQ}(\omega)=\frac{f_{k+(1/2)Q}^*-f_{k-(1/2)Q}^*}{\hbar\omega+\Delta E_{kQ}+i\hbar A_Q}. \tag{29}$$

Summing Eq. (28) over *k* and after rearranging terms we obtain that

$$n(Q,\omega)=\frac{r_0\mathcal{V}(Q)\mathcal{F}(Q,\omega)}{1-\mathcal{V}(Q)\mathcal{F}(Q,\omega)}, \tag{30}$$

where

$$\mathcal{F}(Q,\omega)=\sum_k \mathcal{F}_{kQ}(\omega). \tag{31}$$

Using Eq. (30) in Eq. (7) it follows that

$$\epsilon(Q,\omega)=1-\mathcal{V}(Q)\mathcal{F}(Q,\omega), \tag{32}$$

an expression which for $A_Q=0$ resembles the well-known Lindhart dielectric function; here, however, through this contribution, it depends on the level of excitation created by the pumping source, whose intensity is present in determining T_c^* and A_Q. As we have seen, A_Q is proportional to such an intensity, and T_c^* is determined by Eq. (25), and therefore by E_0. The energy of the holes is composed of two terms, namely, the thermal kinetic energy and the constant energy provided by the source which is storaged in steady-state conditions. The latter is proportional to the intensity of the source, the proportionality coefficient being an absorption coefficient we call G corresponding to the particular kind of excitation to be used (see below). Thus, in general we can write

$$\tfrac{1}{2}nk_BT_c^* = \tfrac{1}{2}nk_BT_0 + GI_0, \tag{33}$$

where T_0 is the temperature of equilibrium with a thermal reservoir in the initial condition of preparation of the sample, and I_0 is, we recall, the uniform intensity of the pumping source.

We notice, for example, that if the pumping source consists of the presence of a constant electric field of intensity $\mathcal{E}$, then in the Ohmic regime we must expect a law of the type $T_c^* = T_0 + \mathcal{D}\mathcal{E}^2$, where $\mathcal{D}$ is a constant proportional to the conductivity and, according to electrodynamics, $\mathcal{E}^2$ is proportional to the energy of the applied field. It can be noticed that in this case the hole's linear momentum must be incorporated in the set of basic variables of Eq. (6), since a current is then present. The situation is similar in the case of incidence of electromagnetic radiation (e.g., from a laser) where we do also have a proportionality with the square modulus of the electric field of the radiation and then—through the Poynting theorem—on the power-flux intensity of the laser beam [10]. In the case of biopolymers under dark excitation it would depend on the metabolic energy-transfer process [5].

We proceed next to an analysis of an optical property, the Raman scattering spectrum, of this system of charged particles in Q1DS out of equilibrium, looking for its dependence on the intensity of the pumping source, that is, on its ''distance'' from equilibrium.

III. RAMAN SCATTERING SPECTRA

The Raman scattering differential cross section is given by

$$\frac{d^2}{d\omega\, d\Omega}\sigma(Q,\omega) = -\mathcal{A}(Q)[1-e^{-\beta^*\hbar\omega}]^{-1}\,\mathrm{Im}\,\epsilon^{-1}(Q,\omega), \tag{34}$$

where Im stands as usual for the imaginary part and $\mathcal{A}(Q)$ is an amplitude [19], and Ω is the solid angle, in stereoradians, presented by the optical window in the detector when looked from the center of the active volume of the sample. In the figures below the amplitude of proportionality (a constant for given Q in the experiment; we recall that $\hbar\omega$ and $\hbar Q$ are the energy and momentum transferred in the scattering event) is ignored, and then the intensity of the Raman signal is given in arbitrary units.

We obtain for the real and imaginary parts of the complex dielectric function of Eq. (32) that

$$\mathrm{Re}\,\epsilon(Q,\omega) = 1 - \mathcal{V}(Q)\sum_k \mathcal{W}_k(Q,\omega)\mathcal{D}_k^{-1}(Q,\omega), \tag{35}$$

$$\mathrm{Im}\,\epsilon(Q,\omega) = \mathcal{V}(Q)\sum_k \hbar A_Q\mathcal{D}_k^{-1}(Q,\omega), \tag{36}$$

where

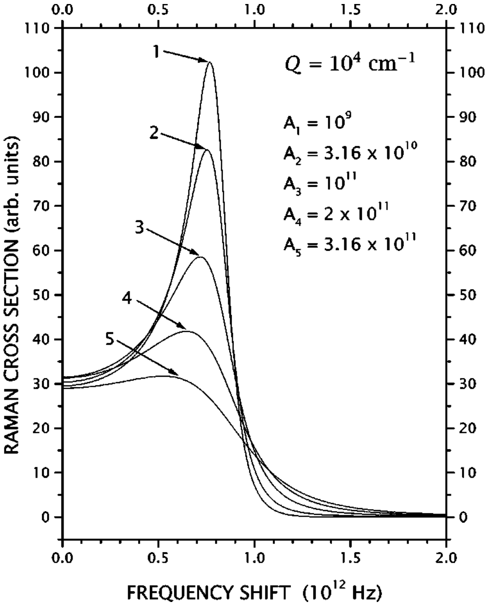

FIG. 1. Raman spectra for different values of the amplitude A of the exciting source and $Q = 0.1\ \mathrm{cm}^{-1}$.

$$\mathcal{W}_k(Q,\omega) = [\hbar\omega + \Delta E_{kQ}][f^*_{k+(1/2)Q} - f^*_{k-(1/2)Q}], \tag{37}$$

$$\mathcal{D}_k(Q,\omega) = [\hbar\omega + \Delta E_{kQ}]^2 + [\hbar A_Q]^2, \tag{38}$$

We proceed next to perform numerical calculations of the Raman spectrum for different levels of excitation in these Q1DS. We use, just for illustration, $R = 250$ Å, $T_c^* = 300$ K, $n_h = 8\times10^5\ \mathrm{cm}^{-1}$, A_Q is left as an open variable parameter, and we use $Q = 0.1\ \mathrm{cm}^{-1}$. In Fig. 1 are displayed spectra corresponding to several different values of A_Q. We can notice that two bands can be identified: they correspond—as in the case of bulk matter—to the one of single-hole excitations (at low-frequency transfer) and to plasma excitations (at higher frequencies). Two relevant results can be pointed out.

(1) With increasing values of the pumping source the band corresponding to scattering by plasma excitations shifts its peak towards lower frequencies. But, more important, it broadens and disappears within the band corresponding to the scattering by the continuum of single-particles excitations. This is a consequence that the contribution containing A_Q in Eq. (16) has the role of a dissipative term, involving the decay of plasma waves in a process mediated by the potential interaction with the external medium, the pumping source in this case, after the collisions with the lattice vibrations have been disregarded. The source then has a twofold role, on the one hand it provides kinetic energy to excite the individual single particles [cf. Eqs. (25) and (33)] and on the other it produces the intense decay of the charge-density wave (the hole plasmons).

(2) The position of the peak of the band corresponding to scattering by plasmons corresponds—as we know from the theory of the Fermi liquid [20]—to the vibrational frequency of the plasma wave. This frequency, in the case of the near one-dimensional system like the one we are considering, can be very approximately derived from a zero of the real part of

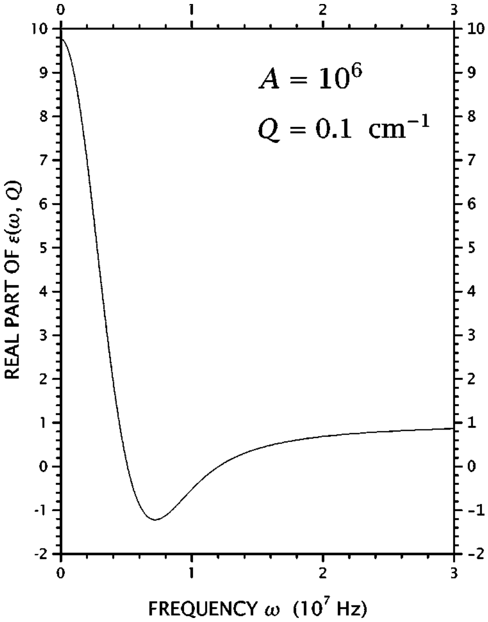

FIG. 2. The real part of the wave number and frequency-dependent dielectric function, for a given value of the amplitude of the exciting source and wave number.

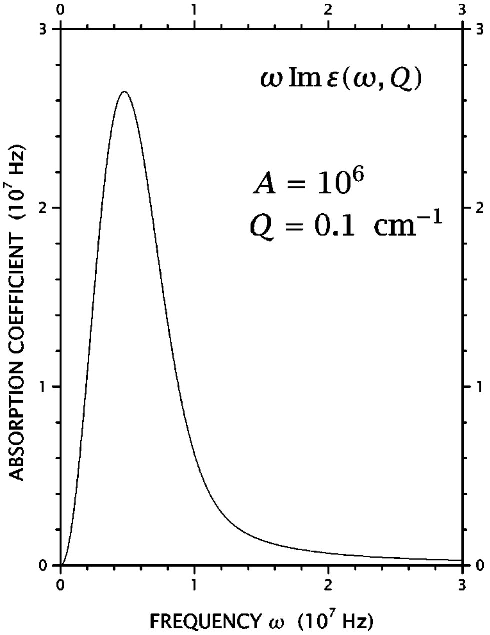

FIG. 3. The absorption coefficient, related to the imaginary part of the dielectric function, for a given value of the amplitude of the exciting source and wave number.

the dielectric function. If we disregard A_Q in the expression for the latter, the plasma-wave dispersion relation is given by [4,21]

$$\omega_Q^2 = \xi Q^2 K_0(QR) \simeq \xi Q^2 |\ln(QR)|, \qquad (39)$$

the approximate value valid for $QR \ll 1$ (the long-wavelength limit), and where ξ is the square of a velocity given by

$$\xi = \frac{2e^2}{\varepsilon_0}\frac{n_h}{m^*}. \qquad (40)$$

The frequency of Eq. (39) is approximately $\omega_Q \simeq 7.5\times 10^{12}\,\mathrm{s}^{-1}$ for the numerical values we are using. This value is in a reasonably good agreement for the lowest value of A in curve 1, while with increasing A there is a shift towards smaller values due to the renormalization that the term in A introduces. To better illustrate the point we consider the case of small A and also small Q (the limit of very long wavelengths), choosing again $Q = 0.1\,\mathrm{cm}^{-1}$ and $A = 10^6$. The real part of the dielectric function is shown in Fig. 2: We can see that it has two zeros; the first one at lower frequencies is associated with single-particle excitations, while the second at higher frequencies is the plasma frequency. The value of the latter agrees very well with the value calculated from Eq. (39). Moreover, the first zero occurs at a frequency given very approximately by $v_{th}Q$, where v_{th} is the thermal velocity $m v_{th}^2 = k_B T_c^*$, as it is also the case for bulk systems.

Finally, in Fig. 3 we show the corresponding optical absorption coefficient defined by [10]

$$\alpha(Q,\omega) = \omega\,\mathrm{Im}\,\epsilon(Q,\omega), \qquad (41)$$

in the conditions of Fig. 2, which shows an absorption band at low frequencies corresponding to transitions between individual single-particle states, and the one due to plasma waves is strongly damped (embedded in the continuum of quasiparticles) and not observable.

We also notice that the frequency and wave-number conductivity $\sigma(\vec{Q},\omega)$ can be derived from the dielectric function through the relation [10]

$$\sigma(Q,\omega) + \frac{i\omega}{4\pi}[\epsilon(Q,\omega) - 1]. \qquad (42)$$

Summarizing, from the Raman scattering spectra the spectrum of excitations of the system in nonequilibrium conditions has been obtained. It consists of the continuum of single-particle excitations (transitions between states in the band of energy levels of holes or the valence band). The other band corresponds to scattering by the collective excitation consisting of the plasma waves. As shown, the energy dispersion relation for these plasmons can be characterized from the optical properties. This is approximately given by Eq. (39), which has a logarithmlike singularity for very small values of wave numbers times the radius of the cylinder-type sample. The group velocity of propagation, at small QR, is

$$v_G(Q) = \frac{d\omega_Q}{dQ} \simeq \xi^{1/2}|\ln(QR)|^{1/2}, \qquad (43)$$

which also has a logarithmlike singularity at the limit of infinite wavelength (of course, such limit is not accessible in real situations with finite samples when the lowest wave-number mode that can be excited is of the order of $Q = \pi/L$).

IV. CONCLUSIONS

We have considered a plasma, composed of holes in geometrically constrained materials (of the ‘‘quantum wire’’ type) as, for example, p doped near one-dimensional semiconductors, polymers, or biopolymers—we recall that in biological systems proteins are p doped [3]. This plasma is under the action of an external pumping source which drives it away (near or far) from equilibrium, through, say, interaction with electric fields—for example, in semiconductors or biological membranes—with electromagnetic radiation such as microwaves, or by the action of the so-called dark excitations (biochemical processes) in biosystems.

The macroscopic state of this out-of-equilibrium system has been described by the thermodynamics of irreversible processes based on a nonequilibrium ensemble formalism, namely, MaxEnt-NESOM. We have concentrated the attention on the optical properties of the electronic system in the plasma, and particularly the influence of the intensity of the source driving the system to a ''distance'' from equilibrium. The frequency and wave-number-dependent dielectric function in such nonequilibrium conditions was derived, since from it one can obtain all the optical properties of the system as well as information on transport properties.

Moreover, the presence of the pumping source (electric field, electromagnetic radiation, dark excitation, or biochemical pumping), drives the holes out of equilibrium, to a thermodynamic state which is characterized by, besides the concentration fixed in the doping process, an excess of kinetic energy out of the equilibrium value characterized by the quasitemperature T_c^*: Its dependence on time reflects the fact that the nonequilibrium state of the system is evolving in time as a result of the irreversible processes that develop in the medium, but, as noticed, when under a continuous and constant source of excitation a time-independent steady state sets in. After the steady state is achieved the nonequilibrium macrostate of the system is described by the statistical operator of Eq. (20). On the other hand, the interaction of the charge-density wave (the plasmons) with the pumping source gives way to a dissipative process [characterized by the contribution $i\hbar A_{kQ}$ in Eq. (16); we recall that compared with it the other contributions arising out of interaction with the lattice vibrations can be disregarded]. This relaxation process is clearly evidenced by the diminution in amplitude and broadening of the band due to scattering by plasma waves. This band dissappears with the increasing intensity of the source (for $A_Q \approx 6\times 10^{11}$ in Fig. 1), and the band due to scattering by single particles is reduced in amplitude.

These are the physical consequences to be expected in near-one-dimensional systems such as semiconductors quantum wires, polymers, and biopolymers, which are driven out of equilibrium by the action of an external source of excitation.

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