Emission of plasmons caused by quasiparticle injection to a high-T_c superconductor

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When the nonequilibrium electron distribution is created by intensive tunneling injection into a cuprate superconductor, the transverse monochromatic Josephson plasma wave is effectively generated in the process of electron energy relaxation. We calculate the generation rate of the Josephson plasmons and find the conditions of coherent electromagnetic emission. [S0163-1829(99)01921-9]

The electromagnetic properties of cuprate superconductors attract significant interest. Due to the high anisotropy of the layered crystal lattice, the ab plane and the c-axis electric transport have a distinct origin. Across the layers, the supercurrent is due to the intrinsic Josephson effect[1] caused by the weak coupling of the adjacent superconducting atomic CuO_2 planes through the interstitial regions. The crystal itself constitutes an array of microscopic Josephson junctions with well established coherence. It results in the appearance of the Josephson plasma mode (JPM) composed by the interlayer electric current and by the c-axis-polarized electromagnetic (EM) field. Since the plasma frequency Ω_{pl} typically is much lower compared to the superconducting energy gap Δ and to the optical phonon frequencies ω_c, most of plasma damping processes are prohibited to occur in the superconducting state.[2] Actually, a sharp resonance peak due to the JPM has been observed.[3] In this paper we show that when quasiparticles are intensively injected into the cuprate superconductor S from the normal metal N through the insulating tunneling barrier I (in so-called N-I-S junction), the Josephson plasma oscillations are excited in S by the recombination of the quasiparticles into the superfluid condensate and due to the elementary electron-plasmon scattering processes. The excited JPM is converted then into the electromagnetic wave in vacuum whose linewidth is determined by the width of the plasma resonance, and thus is extremely narrow, $\delta\Omega \propto \mathrm{Re}\,\sigma_\perp|_{\omega=\Omega_{\mathrm{pl}}}$ ($\sigma_\perp$ is the c-axis ac conductivity). The plasma excitation by quasiparticle injection is discussed in the following way. An intensive injection of quasiparticles causes the deviation of their distribution function f_ε from equilibrium. In the nonequilibrium state f_ε is calculated implementing a kinetic equation. Using the obtained distribution function, and taking into account the Josephson coupling between the CuO_2 planes in the S electrode, we calculate the transverse dielectric function $\epsilon(\omega)$ in the nonequilibrium state. The plasma frequency is determined from the condition that $\mathrm{Re}\{\epsilon(\omega)\}$ vanishes. The complex wave number of the transverse plasma wave inside S is defined as $k=(\omega/c)\sqrt{\epsilon(\omega)}$. Its imaginary part, $k_2=\mathrm{Im}k$, gives the damping constant of the JPM, which is usually small and positive. We show that due to an instability transition, k_2 may become negative for special sets of the electron-plasma collision rate, the tunneling rate, and the bias voltage, indicating the generation of c-axis-polarized coherent plasma waves. The power of the EM emission in the last case is estimated to be quite strong.

The transverse dielectric function is expressed as $\epsilon(\omega) = \epsilon_\infty - \{4\pi i \sigma_\perp(\omega)\}/\omega$, ϵ_∞ being the high-frequency dielectric constant, and $\sigma_\perp(\omega)$ is calculated from the microscopic tunneling approach.[4] To conduct the computations, we assumed (see, also, Ref. 4) that the electron transport is a tunneling along the c axis, while within the ab planes it is metalliclike. The quasi-two-dimensional superconductivity in the planes is described in terms of the Cooper pairing. Then the components of the c-axis ac conductivity in the layered superconducting sample S are written in the general form[4]

$$\sigma_1(\omega) = \frac{\sigma_{N\perp}}{\omega}[\mathcal{I}_{\mathrm{qp}}(\omega)+\mathcal{I}_{\mathrm{Jos},2}(\omega)],$$

$$\sigma_2(\omega) = -\frac{\sigma_{N\perp}}{\omega}[\mathcal{I}_{\mathrm{qp},1}(\omega)+\mathcal{I}_{\mathrm{Jos},1}(\omega)], \tag{1}$$

where $\sigma_1=\mathrm{Re}\,\sigma_\perp$, $\sigma_2=\mathrm{Im}\,\sigma_\perp$, and $\sigma_{N\perp}$ is the average normal-state interlayer tunneling conductivity that in the low-frequency region may be approximated by a constant. The tunneling functions $\mathcal{I}_{\mathrm{qp}}$ and $\mathcal{I}_{\mathrm{qp},1}$ in Eq. (1) are proportional to the contributions from the quasiparticle current (active and reactive, respectively), while $\mathcal{I}_{\mathrm{Jos},1}$ and $\mathcal{I}_{\mathrm{Jos},2}$ are related to the Josephson and interference components of the current.[5] In the nonequilibrium state, the expressions for the functions $\mathcal{I}(\omega)$ can be obtained using the approach[4] where the c-axis electric transport is described as the incoherent tunneling. For the sake of simplicity we neglect the gap anisotropy effects. Then one obtains $\mathcal{I}_{\mathrm{qp}}(\omega)=\int d\varepsilon \tilde{u}_{\varepsilon-\omega}\tilde{u}_\varepsilon F^1_{\varepsilon,\omega}$, $\mathcal{I}_{qp,1}(\omega)=\int d\varepsilon\ \tilde{u}_\varepsilon[\breve{u}_{\varepsilon-\omega}+\breve{u}_{\varepsilon+\omega}]F^2_\varepsilon$, $\mathcal{I}_{\mathrm{Jos},1}(\omega)=\int d\varepsilon\ \tilde{v}_\varepsilon[\breve{v}_{\varepsilon+\omega}+\breve{v}_{\varepsilon+\omega}]F^3_\varepsilon$, $\mathcal{I}_{\mathrm{Jos},2}(\omega)=\int d\varepsilon \tilde{v}_{\varepsilon+\omega}\tilde{v}_\varepsilon F^4_{\varepsilon,\omega}$, $F^1_{\varepsilon,\omega}=n_{\varepsilon-\omega}-n_\varepsilon$, $F^2_\varepsilon=(1-2n_{|\varepsilon|})/2$, $F^3_\varepsilon=(\overline{n}_{|\varepsilon|}-n_{|\varepsilon|})/2$, $F^4_{\varepsilon,\omega}=(n_{\varepsilon-\omega}-\overline{n}_{\varepsilon-\omega}-n_\varepsilon+n_\varepsilon)/2$, that in the equilibrium coincide with the formulas derived in Ref. 4. In our notations, $\tilde{u}_\varepsilon = \mathrm{Re}\, u_\varepsilon$, $\tilde{v}_\varepsilon = \mathrm{Re}\, v_\varepsilon$, $\breve{u}_\varepsilon = \mathrm{Im}\, u_\varepsilon$ and $\breve{v}_\varepsilon = \mathrm{Im}\, v_\varepsilon$, $u_\varepsilon = \varepsilon/\xi_\varepsilon$, $v_\varepsilon = \Delta/\xi_\varepsilon$, $\xi_\varepsilon = \sqrt{\varepsilon^2 - \Delta^2}$.

The distribution functions of the electrons n_ε and of the holes $\overline{n}_\varepsilon$ satisfy $\overline{n}_\varepsilon = 1 - n_\varepsilon$ [in equilibrium, $n_\varepsilon = n_\varepsilon^{(0)} = (\exp(\varepsilon/T)+1)^{-1}$ is the Fermi distribution]. One can also introduce the distribution function of Bogolubov quasiparticles that is connected with n_ε and $\overline{n}_\varepsilon$ as $f_\varepsilon = n_\varepsilon \theta(\varepsilon) + (1 - n_\varepsilon)\theta(-\varepsilon)$ and $f_{-\varepsilon} = \overline{n}_\varepsilon \theta(\varepsilon) + (1 - \overline{n}_\varepsilon)\theta(-\varepsilon)$. From the aforesaid one can infer that the nonequilibrium effect is accounted for the presence of the quasiparticle distribution function f_ε. This function can change its shape under external influence. Here we would like to consider the N-I-S junction where the plasmon emission can be achieved in conditions of a nonequilibrium instability. In this kind of device, the intensive quasiparticle injection from the adjacent N electrode creates the nonequilibrium state in the S electrode. The average electron energy is also increased, and then it dissipates due to inelastic elementary collision processes. The mechanism of the electron energy relaxation strongly depends on the interaction with other excitations in the system. Although in general, the excitation spectrum may be quite complicated, the matter of interest in this paper belongs to a relatively low-energy region $\omega \sim \Omega_{\mathrm{pl}} < \Delta$. Due to the sharp plasma resonance,[2,3] the probability of the electron-plasma collisions is expected to dominate over the probability of other processes. Then, the quasiparticle distribution function is determined from the following kinetic equation:[6,7]

$$\tilde{u}_\varepsilon \dot{f}_\varepsilon = \mathcal{J}^{\text{e-pl}} + \mathcal{J}^{\text{e-ph}} + \tilde{u}_\varepsilon (\dot{f}_\varepsilon)_T, \tag{2}$$

where $\mathcal{J}^{\text{e-pl}}$ is the collision integral of electrons with plasmons

$$\mathcal{J}^{\text{e-pl}}\{f_\varepsilon, N_\omega\} = \alpha_{\mathrm{pl}} \int_{-\infty}^{\infty} d\varepsilon' [\Xi_1 \delta(\varepsilon' - \varepsilon - \Omega_{\mathrm{pl}}) + \Xi_2 \delta(\varepsilon - \varepsilon' - \Omega_{\mathrm{pl}}) + \Xi_3 \delta(\varepsilon' + \varepsilon - \Omega_{\mathrm{pl}})], \tag{3}$$

$$\Xi_1 = (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} + \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \pm 1)[f_{\varepsilon'} - f_{\pm\varepsilon}] + (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} + \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \mp 1) \times [f_{-\varepsilon'} - f_{\pm\varepsilon}], \tag{4}$$

$$\Xi_2 = (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} + \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \pm 1)[f_{\varepsilon'} - f_{\pm\varepsilon}] + (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} + \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \mp 1) \times [f_{-\varepsilon'} - f_{\pm\varepsilon}], \tag{5}$$

$$\Xi_3 = (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} - \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \mp 1)[1 - f_{\varepsilon'} - f_{\pm\varepsilon}] + (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} - \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \pm 1)[1 - f_{-\varepsilon'} - f_{\pm\varepsilon}], \tag{6}$$

where $\alpha_{pl} = (2/3)(l_\perp^2/\tau_\perp)(eE_{\mathrm{pl}}/\Omega_{\mathrm{pl}})^2$ is the electron-plasmon collision rate that must be determined self-consistently, $l_\perp$ is the c-axis electron mean free path, $\tau_\perp$ is the c-axis scattering time, and E_{pl} is the c-axis-polarized electric-field amplitude of the elementary Josephson plasma oscillations. Formally, E_{pl} is related to the interlayer phase difference $\varphi_{n,n-1}$ by the Josephson relationship $E_{\mathrm{pl}} = (\hbar/2e)(\dot{\varphi}_{n,n-1}/c_\perp)$. The δ functions entering Eq. (3) describe the electron energy conservation in the electron-plasmon collision processes, while the terms like $\tilde{u}_\varepsilon \tilde{u}_{\varepsilon'} \pm \tilde{v}_\varepsilon \tilde{v}_{\varepsilon'} \pm 1$ are related to the coherence factors. Contrary to the electron-plasmon interaction, that have a resonant character, the electron-phonon collision term acts in a wide energy range and is approximated here as $\mathcal{J}^{\text{e-ph}}(f_\varepsilon) \approx -(f_\varepsilon - f_\varepsilon^{(0)})/\tau_{\text{e-ph}}$ where $\tau_{\text{e-ph}}$ is the respective relaxation time. The last term in formula (2) is the tunnel source for the N-I-S junction that according to Refs. 6,7 is $u_\varepsilon \cdot (\dot{f}_\varepsilon)_T = \alpha_T Q_T(f_\varepsilon)$, where $\alpha_T = [4\pi^2 e^2 N(0) R_N]^{-1}$ is the tunneling rate,[6] $N(0)$ is the density of electron states at the Fermi level, and R_N is the normal-state resistance of the junction. This kind of source has a property $Q_T(f_\varepsilon) \neq Q_T(f_{-\varepsilon})$ that results in the charge imbalance phenomena[8] and causes the shift of chemical potential (see, also, Ref. 9).

In this paper we implement the above formulas to consider the nonequilibrium emission of the EM wave from the S electrode in conditions of the intensive tunneling injection of quasiparticles using the N-I-S setup. The emission function is obtained from the plasmon-electron collision integral in the form

$$\mathcal{G} = \mathcal{F}(\omega)[\mathcal{G}_1 + \mathcal{G}_2], \tag{7}$$

where $\mathcal{F}(\omega)$ is the plasmon density of states having a sharp resonance maximum at $\omega = \Omega_{\mathrm{pl}}$, $\mathcal{G}_1 = \int_{-\infty}^{\infty} d\varepsilon (\tilde{u}_\varepsilon \tilde{u}_{\omega-\varepsilon} - \tilde{v}_\varepsilon \tilde{v}_{\omega-\varepsilon}) f_\varepsilon f_{\omega-\varepsilon}$ is related to the contribution of elementary recombination processes when, e.g., two electrons are converted into a condensate with emission of a plasmon, and $\mathcal{G}_2 = \int_{-\infty}^{\infty} d\varepsilon (\tilde{u}_\varepsilon \tilde{u}_{\varepsilon-\omega} + \tilde{v}_\varepsilon \tilde{v}_{\varepsilon-\omega}) f_\varepsilon (1 - f_{\varepsilon-\omega})$ corresponds to the relaxation part when the electron loses its energy emitting a plasmon. The above Eqs. (2)–(7) must be completed by the self-consistency equation for the superconducting order parameter that also is affected in the nonequilibrium state.[10] For the sake of simplicity, we describe the SC order parameter in terms of the effective temperature $T^\star$, that is used here as a parameter.

At the beginning, we apply Eqs. (1) to calculate the equilibrium dielectric function $\epsilon(\omega)$ that depends on the Fermi distribution function of electrons $n_\varepsilon^{(0)} = 1/[\exp(\varepsilon/T)+1]$. For certainty we accept the following normal-state parameters: $\epsilon_\infty = 23$, $\sigma_{N\perp} = 1.6\ \Omega^{-1}\ \mathrm{cm}^{-1}$. The calculated results for the real part $\epsilon_1(\omega)$ are presented in Fig. 1. The curve A in Fig. 1 belongs to the equilibrium case at temperature $T = 0$ with no bias voltage applied across the junction, $V = 0$. One can see that $\epsilon_1(\omega)$ vanishes at $\omega = 0.74$ (the frequency ω, the bias voltage V, the electron energy ε and the temperature are expressed in units of the energy gap magnitude Δ at $T = 0$). At this frequency $\omega = 0.74$, one obtains $\epsilon_2 = \mathrm{Im}\{\epsilon(\omega = 0.74)\} \approx 0.5$ that is small enough to provide a sharp resonance (JPM) with $\Omega_{\mathrm{pl}}^{(0)} = 0.74$ [case (i)]. The magnitude of the plasma frequency in cuprate SC $\Omega_{\mathrm{pl}} \simeq 0.5 - 10$ mV is typically smaller compared to the value of the energy gap $\Delta \approx 10 - 40$ mV.

To find the quasiparticle distribution function in the nonequilibrium but stationary case $\dot{f}_\varepsilon = 0$, we reduce Eq. (2) to a set of linear differential equations with respect to functions of the kind $f_{\pm\varepsilon}$, $f_{\pm\varepsilon\pm\Omega_{\mathrm{pl}}}$, $f_{\pm\varepsilon\pm V}$ and their derivatives. This is achieved, e.g., if one implements the condition that at

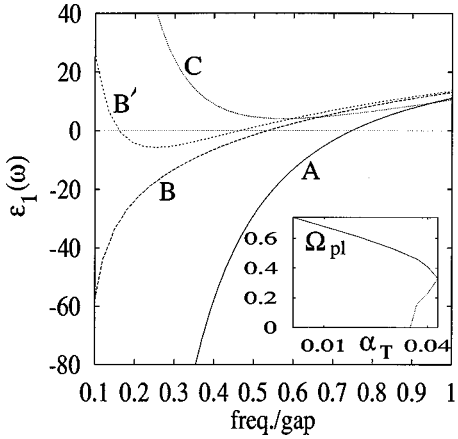

FIG. 1. The real part of the dielectric function $\epsilon_1(\omega)$ in the layered superconductor for different parameters of the quasiparticle injection and of the electron energy relaxation. The inset shows the dependence $\Omega_{\rm pl}(\alpha_T)$ for the same parameters as for curves B and B′.

large ε (namely $\varepsilon \gg \max\{\Delta,\Omega_{\rm pl},V\}$) $f_\varepsilon \Rightarrow f_\varepsilon^{(0)}$. Besides we assume that the electron-phonon collisions are weak, because the peaks in the phonon density of states are situated at higher frequencies. At this point, we have to specify the nonequilibrium value of the plasma frequency $\Omega_{\rm pl}$ that is present in Eq. (3), and that must be determined self-consistently. To compute $\Omega_{\rm pl}$, we implemented the following recursion procedure. Initially we accepted the trial value of the plasma frequency $\Omega_{\rm pl}=\Omega_{\rm pl}^{(0)}$. At a fixed bias voltage V, for given rates of the tunneling injection α_T and of the electron-phonon collision $\alpha_{\text{e-ph}}=\hbar/\tau_{\text{e-ph}}$, we solved the kinetic equation (2) and obtained the nonequilibrium function f_ε. Then, using Eqs. (1), we calculate the dielectric function $\epsilon(\omega)$. Implementing the condition of the plasma resonance $\epsilon_1(\omega)=\text{Re}\{\epsilon(\omega)\}|_{\omega=\Omega_{\rm pl}}=0$, one obtains the nonequilibrium value of the plasma frequency $\Omega_{\rm pl}^{(1)}$ in the first iteration. According to Eqs. (1), $\epsilon(\omega)$ depends on the shape of the distribution function, therefore the position of the plasma resonance $\Omega_{\rm pl}^{(1)}$ is generally shifted with respect to $\Omega_{\rm pl}^{(0)}$. Then, the calculated value $\Omega_{\rm pl}^{(1)}$ is used to proceed to the second step, solving the kinetic equation (2) and computing $\epsilon(\omega)$ again to find $\Omega_{\rm pl}^{(2)}$, etc. The procedure is repeated i times until the ratio $|\Omega_{\rm pl}^{(i)}-\Omega_{\rm pl}^{(i-1)}|/\Omega_{\rm pl}^{(i)}$ becomes sufficiently small, giving $\Omega_{\rm pl}^{(i)} \to \Omega_{\rm pl}$. The nonequilibrium alternation of $\Omega_{\rm pl}$ is well illustrated by Fig. 1 where we plot $\epsilon_1(\omega)$ for a different intensity of the quasiparticle injection and of the electron energy relaxation rates. At finite bias voltage, the quasiparticle tunneling injection from the N-electrode changes the shape of f_ε in S, causing the shift of JPM. This tendency is observed already at a relatively small intensity of the quasiparticle injection. From curve B of Fig. 1 [case (ii)] corresponding to $V=4.95$, $T^\star=0.3$, $\alpha_T=0.03$, $\alpha_{\text{e-pl}}=2.15$, and $\alpha_{\text{e-ph}}=0.19$ one finds the new value $\Omega_{\rm pl}=0.53$ that is significantly smaller compared to $\Omega_{\rm pl}^{(0)}$. At some level of tunneling rate α_T, the nonequilibrium dielectric function can vanish twice, at $\Omega_{\rm pl}^a$ and $\Omega_{\rm pl}^b$. That can be seen, e.g., from curve B′ calculated for the same parameters as before but with $\alpha_T=0.037$. It gives two Josephson plasma frequencies, $\Omega_{\rm pl}^a \simeq 0.17$ and $\Omega_{\rm pl}^b \simeq 0.46$. From the inset to Fig. 1 one can also infer that the

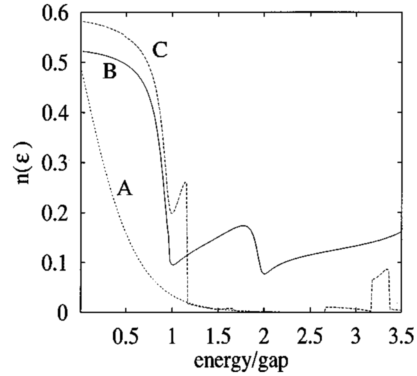

FIG. 2. The nonequilibrium distribution functions of electrons n_ε for the same parameters of injection and relaxation as in the previous figure.

dependence $\Omega_{\rm pl}$ versus α_T (all the other parameters are the same as before) is essentially nonmonotonic with the two branches at $\alpha_T > 0.036$. If the intensity of the quasiparticle injection is increased, the JPM may be even washed out. One can see that from curve C [where we use $V=2.2$, $T^\star=0.3$, $\alpha_T=0.91$, $\alpha_{\text{e-pl}}=0.1$, and $\alpha_{\text{e-ph}}=5.3$, case (iii)] of the same Fig. 1 where $\epsilon_1(\omega)>0$ in the whole frequency range and diverges at $\omega \to 0$. The small value of $\alpha_{\text{e-pl}}$ and relatively large value of $\alpha_{\text{e-ph}}$ is used because the electron energy relaxation mechanism for the case (iii) is mostly due to the nonresonant electron-plasmon and electron-phonon collisions (just like in classic superconductors). From this curve C one can see that the JPM is entirely vanished in such conditions. In this limit, the calculations show that the imaginary part of the dielectric function $\epsilon_2(\omega)|_{\omega\to 0} \Rightarrow 0$ becomes small at the low frequency $\omega \to 0$, providing thus a transparency window for EM waves. The disappearance of JPM corresponds to a nonequilibrium instability of S with respect to its screening properties, causing a static homogeneous electronic ferroelectric state.

If the tunneling injection from the N electrode to S is perpendicular to the atomic CuO_2 planes, the nonequilibrium distribution[6,7,10] is established on the diffusion length $l_D \simeq \sqrt{D\tau_\varepsilon}$, where τ_ε is the electron energy relaxation time, $D \simeq l_{\rm el} v_F/3$, v_F being the Fermi velocity, although the electron elastic mean free path $l_{\rm el}$ along the c axis in cuprates is quite small[11] ($l_{\rm el} \sim c_\perp \sim 1$ nm). For a cuprate SC, $l_D/l_{\rm el} \simeq \sqrt{\tau_\varepsilon/3\tau_{\rm el}} \approx 10-10^2$ where $\tau_{\rm el}$ is the elastic scattering time. Thus, the nonequilibrium region may involve at least $\sim 10-10^2$ layers. For the tunneling injection along the ab planes, the size of the mentioned region can be much larger.

The electron distribution functions n_ε obtained for the nonequilibrium cases (ii-iii) are presented in Fig. 2. The parameters used for calculations of curves A, B, and C are the same as for corresponding curves of Fig. 1. For instance, curve A corresponds to the equilibrium distribution, while curve B in this Fig. 2 is related to n_ε in the limiting case (ii) of relatively low injection.

Due to the elementary electron-plasma interaction, the part of the energy from the injected quasiparticles may be emitted as an EM wave polarized along the c axis. In equilibrium, due to the electron-plasmon and the electron-phonon collisions, there is a uniform background of the emission. At finite bias voltage $V \neq 0$, the excess energy of injected quasiparticles is consequently transferred to the plasma excita-

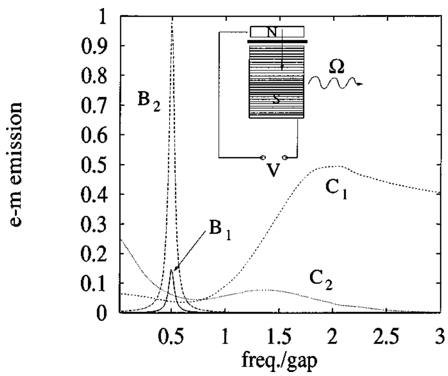

FIG. 3. The emission function of far infrared.

tions. Because the sharp peak in the density of states of the plasmons $\mathcal{F}(\omega)$ in the S electrode, the emission $\mathcal{G}(\omega)$ has a maximum at $\omega=\Omega_{\rm pl}$. However, if the tunneling injection from the adjacent N electrode is weak, the emitted power remains insignificant. As the level of injection grows, it increases the external energy supply to the cuprate SC electrode, providing the monochromatic emission as shown by curves B_1 and B_2 [in this case (ii), B_1 is related to the function $\mathcal{G}_1$ in Eq. (7), while B_2 corresponds to $\mathcal{G}_2$] in Fig. 3. The width of the maximum in the emission spectrum is remarkably narrow, $\delta\Omega \propto \mathrm{Re}\ \sigma_\perp|_{\omega=\Omega_{\rm pl}}$. As the bias voltage increases, the sharp line in the emission spectrum is gradually transformed into the thresholdlike behavior with a wide maximum at $\omega \simeq 2\Delta$ becoming nonresonant [see curves C_1, C_2 belonging to case (iii)]. The mentioned curves C_1 and C_2 show a finite emission at $\omega \to 0$, corresponding to a tail from the JPM that is strongly damped in case (iii).

Although the emission of plasmons in the above-considered cases (ii)–(iii) is not generally coherent, one may achieve, in principle, a strong coherent EM irradiation from the N-I-S setup as well. Such kind of far-infrared plasma laser may be obtained at the following conditions. For certain shapes of the electron distribution function n_ε, when the decrement $k_2=(\omega/c)\mathrm{Im}\{\sqrt{\epsilon(\omega)}\}$ changes its sign, becoming negative, there occurs another kind of nonequilibrium instability in the electron system of S. It corresponds to the amplifying of the EM wave at the given frequency ω. If the plasma resonance exists, its damping is small.[2] Then any weak disturbance of the JPM is amplified due to the transfer of energy from the nonequilibrium electrons to the resonant plasmons. The phase coherence is established by the initial excitation that develops into the electromagnetic wave. This situation may be, e.g., realized for deviations of kind $\delta f_\varepsilon = (f_\varepsilon - f_\varepsilon^{(0)})/f_\varepsilon^{(0)} \propto e^{-(\omega-\omega_0)^2/\kappa}$ where parameters ω_0 and κ are estimated from the approximate solution of the kinetic equation (2) referring to a definite experimental setup. One can appraise the power generated in the elementary electron-plasma collisions from the balance equation as $P_{\text{e-pl}} = P_T - P_{\text{e-ph}}$, where the power supplied by the tunneling junction is $P_T = \alpha_T \int d\varepsilon\, \varepsilon Q_T(f_\varepsilon)$, while the power scattered due to the electron-phonon collisions is $P_{\text{e-ph}} = \int d\varepsilon\, \varepsilon \mathcal{J}^{\text{e-ph}}(f_\varepsilon)$. If one implements the parameters $V=0.5\cdot\Delta/e \simeq 15$ mV (e.g., for Bi-based compounds), $R_N=10\ \mathrm{k}\Omega$, $N(0)=10^{27}\ \mathrm{m}^3/\mathrm{eV}$, it gives $\alpha_T \simeq \alpha_{\rm pl} \simeq 2\times10^{-5}$, $\kappa \approx 0.3$, and $\omega_0 \simeq 0.3$. Accepting also the values $l_\perp/c_\perp = 5$, $eV_{\rm pl}/\hbar\Omega_{\rm pl} = 0.02$, where $V_{\rm pl} = E_{\rm pl} c_\perp$, $\hbar/\tau_\perp = 0.3$, and $c_\perp = 1$ nm is the c-axis lattice constant, one finds $E_{\rm pl} \approx 15$ V/mm, and the flow of energy per unit area is $P_{\text{e-pl}} \simeq \epsilon_0 \bar{c} E_{\rm pl}^2 \approx 2\ \mathrm{mW/mm}^2$, where $\bar{c} \simeq 10^6$ m/s. Indeed, the generation of the coherent plasma wave takes place if the condition $k_2 < 0$ is fulfilled. The sufficiently strong EM emission was recently observed in the tunneling injection experiments.[12]

Let us emphasize right now the essential difference between the EM emission from the nonequilibrium state of the cuprate superconductor compared to the known effect of the phonon emission from the ordinary low temperature SC. First, the JPM frequency, and thus the frequency of the EM emission $\simeq \Omega_{\rm pl}$ is typically $10-10^2$ times lower compared to the magnitude of the energy gap in cuprates. Contrary, the peak of the phonon emission in ordinary SC is attributed rather to the BCS singularity in the electron density of states at much larger frequencies $\omega \simeq 2\Delta$. Second, the width of the line $\delta\Omega$ is much narrower in the case of the plasmon emission, $\Omega_{\rm pl} \approx \mathrm{Re}\ \sigma_\perp(\omega)|_{\omega=\Omega_{\rm pl}}$, compared to classic SC where the linewidth is determined by the rate of inelastic collisions, i.e., $\delta\Omega \simeq \mathrm{Im}\,\Delta$. Third, the phonon emission from ordinary SC is always incoherent, while from the cuprate electrode, one may achieve the coherent EM irradiation. Finally, the plasma frequency in classic SC materials, $\omega_p \sim 10^{15}\ \mathrm{s}^{-1}$, is several orders of magnitude higher compared to the JPM frequency, $\Omega_{\rm pl} \sim 10^{10} - 10^{12}\ \mathrm{s}^{-1}$. It means that in ordinary SC, the EM field with frequencies below $\omega_{\rm pl}$ may propagate only at surface or in thin films. Contrary, the c-axis-polarized EM waves with frequencies $\omega \leqslant \Omega_{\rm pl}$ may be generated in the bulk of the cuprate SC electrode.

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