Interplay between collective modes in hybrid electron gas–superconductor structures

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We study hybridization of collective plasmon and Carlson-Goldman-Artemenko-Volkov modes in a hybrid system, consisting of a two-dimensional layers of electron gas in the normal state and superconductor, coupled by long-range Coulomb forces. The interaction between these collective modes is not possible in a regular single-layer two-dimensional system since they exist in non-overlapping domains of dimensionless parameter \( \omega \tau \), where \( \omega \) is the external electromagnetic field frequency and \( \tau \) is electron scattering time. Thus, in a single-layer structure, these modes are mutually exclusive. However, the coupling may become possible in a hybrid system consisting of two separated in space materials with different properties, in particular, the electron scattering time. We investigate the electromagnetic power absorption by the hybrid system and reveal the conditions necessary for the hybridization of collective modes.

I. INTRODUCTION

Mesoscopic systems in normal (not superconducting) state lodge elementary excitations, such as excitons and plasmons. They determine the response of the material to external perturbations. Plasmons represent collective modes of particle density fluctuations and play an essential role in the material response to external alternating electromagnetic (EM) fields. Their dispersion law and damping strongly depend on the dimensionality of the system. Indeed, in the case of three-dimensional (3D) electron plasma, the plasmon branch is with a good accuracy dispersionless, whereas two-dimensional (2D) plasmons are characterized by a square-root dispersion relation and vanish in long-wavelength limit (at least, in the framework of the quasistatic approximation). Many-component 2D systems like an electron-hole 2D gas are characterized by two plasmon branches with linear (acoustic plasmon) and square-root (optical plasmon) dispersions.

If with a decrease of temperature the system undergoes a phase transition to another state like a Bose-Einstein or superconducting (SC) condensate, the ground state changes and new elementary excitations describing low-energy properties of the system arise. This can lead to drastic modification of the material response to external EM fields, which becomes especially important in view of recent progress in plasmonics, optoelectronics, and high-temperature condensation phenomena.

In the pioneering work of Bardeen Cooper and Schiffer [1], only single-particle excitations characterized by the SC gap were studied. Later, Anderson [2], Bogoliubov [3] and Vaks, Galitskii, and Larkin [4] developed a more extensive theory, which includes the collective excitations of Cooper pairs. They showed, that in a neutral Fermi system with an attracting interaction between fermions, collective modes possess a soundlike dispersion. A charged Fermi system, instead, demonstrates the collective modes with the frequency pushed towards the plasmon frequency of the 3D charged electron gas. This frequency is usually much larger than the BSC gap and consequently it does not play role in conventional phenomena in SCs.

These modes correspond to oscillations of the SC order parameter. More precisely, the Anderson-Bogoliubov mode represents the oscillation of the phase of the order parameter, whereas the Higgs mode corresponds to the oscillations of its amplitude. These collective modes play crucial role in the gauge-invariant response of SCs to external EM perturbations [3].

Later on, Carlson and Goldman [6,8] experimentally discovered another type of collective modes, which occur in the vicinity of the transition temperature \( T_c \), in SCs. This experimental finding stimulated theoretical studies. In the dirty-sample case, there was developed a model of Schmid and Schön [9], whereas in the clean (low-disorder) limit, there appeared a model of Artemenko and Volkov (AV) [10,11]. Qualitatively, AV mode corresponds to the out-of-phase oscillations of normal and SC currents in such a way that it prevents the emergence of net charges in the system. The frequency of this mode is smaller than the plasmon frequency of normal 3D electrons, thus suppressing the interaction between them.

In the meanwhile, the coupling between single-particle and collective modes in hybrid systems consisting of a 2D electron layer and a 2D condensate of Bose particles represent actively developing areas of research recently [12-14]. In particular, there emerge such phenomena as the magnetoplasmon Fano resonances in uniform magnetic fields [15] and amplification of the incident light [16,17]. The progress in theoretical proposals and the lack of well-established experimental platforms motivates the study...
of the interplay of different collective modes in hybrid systems. In this article, we investigate the hybridization of collective plasmon and AV modes in a hybrid semiconductor-SC structure (where one of the subsystems is in the SC state). As we have already pointed out, the plasmon mode has a gapless dispersion in 2D systems (in the SC state). As we have already pointed out, the plasmon mode has a gapless dispersion in 2D case. We will show that it opens a possibility of interaction between AV modes and plasmons.

At first sight, despite the gaplessness of the plasmon mode, the interaction between them is not possible since the AV mode exists in the domain \( \omega \tau \ll 1 \) (here \( \tau \) is electron-impurity scattering time), whereas the plasmon mode exist at \( \omega \tau \gg 1 \), and these two conditions do not overlap. However, their interaction can become possible if the normal electron gas layer (having the plasmon mode) and the SC layer (hosting the AV mode) are separated in space and electrons in each layer have different scattering times. Thus, if we denote as \( \tau_n \) the scattering time of electrons in the normal layer and \( \tau_s \) the one for quasiparticles in the SC layer, the simultaneous existence of AV and plasmon modes is possible when \( \tau_n^{-1} > \omega > \tau_s^{-1} \). This is a necessary criterion for the coexistence of both the modes in one sample. However, it is not a sufficient condition for their hybridization. What are the other conditions? This is the main question, which we address in this article.

**II. LONGITUDINAL DIELECTRIC FUNCTION AND THE DISPERSION RELATION**

Let us consider a 2D layer of a semiconductor containing normal electrons with the equilibrium density \( N_s \) in the vicinity of a 2D SC layer containing both normal single-electron excitations above the SC gap \( \Delta(T) \), where \( T \) is temperature, and SC Cooper pairs described by the order parameter (Fig. 1). The total electron density in SC layer is \( N_s \), whereas the the density of SC electrons we denote as \( n_s \).

Furthermore, we expose the system to an external EM field, which represents a plane wave with 3D wave vector \( \mathbf{Q} = (k_x, Q_y) \), where the 2D wave vector \( \mathbf{k} = (Q_x, Q_y) \) is its in-plane projection. The field creates electric currents in both the layers,

\[
\begin{align*}
\mathbf{j}_s &= \sigma_s (\mathbf{E}_0 + \mathbf{E}_s'), \\
\mathbf{j}_n &= \sigma_n (\mathbf{E}_0 + \mathbf{E}_n'),
\end{align*}
\]

where\( \mathbf{E}_0 \) is the in-plane component of external EM field, which we assume to be equal in both layers (valid as long as \( Q_x d \ll 1 \)); \( \mathbf{E}_{n,s}' \) are the internal EM fields, induced in the corresponding layers due to the oscillations of the particle densities;

\[
\sigma_n = \frac{e^2 N_s \tau_n / m_n}{1 - i\omega \tau_n}
\]

is a conventional dynamical Drude conductivity of the degenerate electron gas in the semiconductor, with \( m_n \) being the electron effective mass; \( \sigma_s \) describes the current response of single-particle excitations in the SC layer. It should be noted, that in Eq. (1) we omitted the factor \( \exp(ikr - i\omega t) \) since all the time and in-plane position-dependent quantities are proportional to this factor.

The induced fields depend on \( z \), \( \mathbf{E}_{n,s}'(z = \pm d) \), where \( \mathbf{E}' \) can be expressed via the scalar potential in the framework of quasistatic approximation, \( \mathbf{E}' = -\nabla \varphi' \). The scalar potential, in turn, satisfies the Poisson equation,

\[
\left( \frac{d^2}{dz^2} - k^2 \right) \varphi'(z) = -\frac{4\pi}{\epsilon} [\delta \rho_s \delta(z - d) + \delta \rho_n \delta(z + d)],
\]

where \( \delta \rho_{n,s} \) are deviations of charge densities from their equilibrium values, and \( \epsilon \) is a dielectric function of the material. Combining (1) and (3) with the continuity equation, we find

\[
\left( \frac{1 + 2\pi i k}{\epsilon \omega} \sigma_s \right) \left( \frac{2\pi i k e^{-2\pi k d \sigma_s}}{\epsilon \omega} \sigma_s \right) \left( \delta \rho_s \right) \frac{k \mathbf{E}_0}{\omega} \left( \sigma_s \right) = \frac{\left( \sigma_s \right)}{\epsilon(k, \omega)}
\]

and

\[
\left( \mathbf{j}_s \right) = \left( \frac{1 + 2\pi i k}{\epsilon \omega} \left( 1 - e^{-2\pi k d \sigma_s} \sigma_s \right) \right) \left( \sigma_s \right) \frac{\mathbf{E}_0}{\epsilon(k, \omega)},
\]

where we have introduced the longitudinal dielectric function,

\[
\epsilon(k, \omega) = \left( 1 + \frac{2\pi i k}{\epsilon \omega} \sigma_s \right) \left( 1 + \frac{2\pi i k}{\epsilon \omega} \sigma_n \right) + \left( \frac{2\pi k e^{-2\pi k d \sigma_s}}{\epsilon \omega} \right)^2 \sigma_s \sigma_n.
\]
Here and in what follows we assume that $k$ and $E_0$ only have $x$ in-plane components.

Eq. (6) is similar in form to the standard one, describing generic two-component systems. Equation $\varepsilon(k, \omega) = 0$ determines the dispersion and damping of collective modes. The key quantity here is $\sigma_s$, which has to be found accounting for the interaction of single-particle excitations in SC layer with the Cooper pair condensate.

Let us, first, switch off the interaction between the layers and analyze the collective modes in each layer separately. Formally, the noninteracting case corresponds to putting $d \to \infty$ in Eq. (6). Then $\varepsilon(k, \omega) = 0$ splits into two independent conditions,

$$1 + \frac{2\pi ik}{\epsilon \omega} \sigma_n = 0; \quad 1 + \frac{2\pi ik}{\epsilon \omega} \sigma_s = 0. \quad (7)$$

Let us consider these cases separately.

A. Plasmons in the layer containing 2DEG

Substituting Eq. (2) in (7), we find the plasmon dispersion and its damping due to the electron-impurity scattering,

$$\omega = \omega_n \sqrt{1 - \frac{1}{(2\omega_n \tau_n)^2}} - i \frac{\pi k}{2 \tau_n}; \quad \omega_n = \frac{2\pi e^2 k N_s}{c m_n}. \quad (8)$$

The plasmon exists if $2\omega_n \tau_n > 1$. In the case of weak scattering, when $\omega_n \tau_n \gg 1$, the plasmon damping is much smaller in comparison with its frequency, and in the EM power absorption spectrum, the plasmons are seen as well-resolved resonances. The spatial dispersion of Drude conductivity (2) does not play role in (8) and can be neglected if $\omega \gg k v_n$, where $v_n$ is the Fermi velocity of electrons in the semiconducting layer.

B. Artemenko-Volkov modes in the SC layer

To find the dispersion law of the AV mode, we have to, first, find the conductivity $\sigma_s$ of single-particle excitations above the SC gap. In their original work [10], Artemenko and Volkov used the quasi-classical approach based on the kinetic equations. Later, the validity of their results was confirmed by the quantum field theory methods [13]. We will thus follow the simpler original Boltzmann equation approach, where the AV modes are found by analyzing the longitudinal dielectric function (7). The calculation of $\sigma_s$ for a 3D SC can be found elsewhere [19], and we adopt it for 2D SC layer just presenting here the result,

$$\sigma_s = \sigma_{0s} \frac{\omega^2 - u^2 k^2 + i \omega \eta_s + i k^2 v^2 \omega \tau_s J_{\omega}}{\omega^2 - u^2 k^2 + i k^2 \omega \tau_s J_{\omega}}, \quad (9)$$

where

$$\sigma_{0s} = \frac{e^2 N_s \tau_s}{m_s}, \quad u = v_s \sqrt{7\zeta(3) \frac{\Delta}{2\pi^3 \omega^4 T}}, \quad \eta_s = \frac{n_s}{N_s} \frac{1}{\tau_s} \quad (10)$$

$$J_{\omega} = \frac{\ln(1/\omega \tau_s)}{\pi}.$$  

Here we introduced a static Drude conductivity of normal electrons in SC layer, $\sigma_{0s}$; $v_s$ is the Fermi velocity in SC layer; $u$ is the phase velocity of AV mode (see the description below), $\zeta(3)$ is the Riemann zeta-function. We note also that $u \ll v_s$.

Formula (9) is valid under the conditions $\omega \gg k v_s, \tau_s^{-1} \ll \Delta \ll T$ and $(\tau_s k v_s)^2 \ll \omega \tau_s \ll 1$ [19]. Using these assumptions and combining Eqs. (9) and (7), we find

$$1 + i \frac{\omega^2 \tau_s}{\omega} \left(1 + \frac{i \omega \eta_s}{\omega^2 - u^2 k^2 + i k^2 v^2 \omega \tau_s J_{\omega}}\right) = 0, \quad (11)$$

where $\omega_s = \sqrt{\frac{2\pi e^2 k N_s}{c m_s}}$.

In the limiting case $\omega_s^2 \tau_s \gg \omega$ [19], there exists an analytical solution,

$$\omega^2 = u^2 k^2 - i k^2 v_s^2 \omega \tau_s J_{\omega} - i \omega \eta_s. \quad (12)$$

We note again, that this solution describes a weakly-damped sound-like AV mode in the frequency range

$$\left(\frac{\Delta}{T}\right)^2 \ll \omega \tau_s \ll \frac{\Delta}{T} \ll 1. \quad (13)$$

The numerical solution of Eq. (11) gives two collective branches in the isolated 2D SC layer. The first dispersion is presented in Fig. 2, where we use dimensionless frequency $\Omega = \omega \tau_s$ and wave vector $q = (\omega, \tau_s)^2$. We see that the AV mode characterized by the linear dispersion $uk$ lies above the almost horizontal part of red-dotted line, which stands for the lower bound $(\Delta/T)^2$ in Eq. (13).

The second mode is located in the low-$q$ domain, as shown in Inset in Fig. 2. At $q \ll 1$, this mode is weakly-damped due to the smallness of its imaginary part ($\sim q$) in comparison with its eigen frequency $\sim \sqrt{q}$. The dispersion of this mode can be found analytically from Eq. (11). Indeed, in the limit $k \to 0$, Eq. (11) reduces to

$$1 + i \frac{\omega^2 \tau_s}{\omega} \left(1 + \frac{i \eta_s}{\omega^2 - u^2 k^2 + i k^2 v^2 \omega \tau_s J_{\omega}}\right) = 0, \quad (14)$$

which gives the dispersion of the 2D plasmon,

$$\omega_s^2 \sqrt{\frac{n_s}{N_s} \left(1 - \frac{N_s}{n_s} \frac{\omega_s \tau_s}{2} - i \omega \tau_s \frac{2}{2}\right)}, \quad (15)$$

existing when $n_s/N_s \gg (\omega_s \tau_s)^2/4$. On condition $n_s/N_s \gg (\omega_s \tau_s)^2/4$, formula (15) turns into not damped 2D plasmon oscillations of superconducting electrons

$$\omega_s \sqrt{\frac{n_s}{N_s} \left(1 - \frac{N_s}{n_s} \frac{\omega_s \tau_s}{2} - i \omega \tau_s \frac{2}{2}\right)}, \quad (16)$$
It should be noted, that in contrast with the normal-state systems, where the plasmon exists if \( \omega_n \tau_n \gg 1 \), the mode Eq. (16) exists at \( \omega_n \tau_n \ll 1 \). This mode reminds the Anderson-Bogoliubov mode [2, 3] studied at \( T = 0 \) in 3D SCs. It represents plasmon oscillations of SC electrons, which density at \( T = 0 \) equals to the total electron density in the SC. The difference is that in our 2D case, at \( T \neq 0 \) this mode is determined by the density of SC electrons \( n_s \neq N_s \). We will call it quasi-Anderson-Bogoliubov mode (qAB). The qAB oscillations of the SC condensate at \( T \neq 0 \) are accompanied by the appearance of induced charges in the system which, obviously, interact with normal electrons. They exist in SC at nonzero temperatures and play the role of a friction force influencing the oscillations of the condensate. At \( T_c - T < T_c \) (the case which we consider in this article), the density of normal electrons in the SC layer is of the order of total electron density \( N_s \). It explains why the imaginary part in Eq. (15) is proportional to \( \omega_s^2 \propto N_s \).

III. INTERACTION BETWEEN THE MODES

We have already found, that the plasmon mode of SC layer lies above the boundary of the particle-hole continuum, \( \omega_n \gg k v_n \), whereas the AV mode has the phase velocity \( u \ll v_s \). Hence there exist two limiting cases, depending on the ratio between the Fermi velocities in SC, \( v_s \), and normal 2DEG, \( v_n \), layers (Fig. 3). Let us consider both of these scenarios independently.

A. Modes interaction if \( v_n > v_s \)

In this case, the dispersions of the modes do not intersect. Thus at a given value of the wave vector \( k \), these modes are excited independently in the system at different frequencies of external EM field. Nevertheless, the presence of the other layer results in the inter-layer Coulomb interaction of the particles and produces an additional damping of the collective mode which is excited. Let us analyse this damping.

In the frequency range when the plasmon mode exists, \( \omega \sim \omega_n, \omega_n \tau_n \gg 1 \) we can simplify formulas for the conductivities (2) and (9),

\[
\sigma_n \approx \frac{e^2 N_n}{-i \omega m_n}, \quad \sigma_s \approx \sigma_{s0} = \frac{e^2 N_s \tau_s}{m_s},
\]

and the dispersion equation \( \varepsilon(k, \omega) = 0 \) reduces to

\[
\left( 1 - \frac{\omega_n^2}{\omega^2} \right) \left( 1 + \frac{\omega_s^2 \tau_s}{\omega} \right) = -i \omega_s \frac{\omega_n^2 \omega_s}{\omega^2} e^{-4kd}.
\]

Its iterative solution gives the renormalization of the plasmon dispersion and an additional damping,

\[
\omega = \omega_n \left( 1 - \omega_s \tau_s \frac{\omega_s \tau_s (1 + e^{-4kd})}{4N_n/N_s + \omega_s^2 \tau_s^2 (1 + e^{-4kd})^2} e^{-4kd} \right).
\]

We see that if \( kd \) is not very large and \( N_n/N_s \ll \omega_s^2 \tau_s^2 \), the presence of SC layer can drastically change the plasmon dispersion resulting in a dramatic red shift of the plasmon frequency.

If, instead, the frequency lies in the range of the AV mode, \( \omega \approx u k \), the Drude conductivity of electrons in the semiconducting layer can be replaced by its static limit
σ_n ≈ ε_n² τ_n/m_n, and the dispersion equation reads

\[\left(1 + i \frac{\omega_n^2 \tau_n}{\omega}\right) \left(1 + i \frac{\omega_s^2 \tau_s}{\omega}\right)(1 - i \frac{\omega \eta_s}{\omega^2} + i k^2 v_d^2 \omega \tau_s f_\omega) \right) = \left(\frac{\omega_n^2 \tau_n}{\omega}\right) \left(\frac{\omega_s^2 \tau_s}{\omega}\right) \times \left(1 + i \frac{\omega \eta_s}{\omega^2} - i k^2 v_d^2 \omega \tau_s f_\omega \right).\]

The bare AV mode can exist under the condition \(\omega_n^2 \tau_n \gg \omega\), as we have discussed above. Due to the relation \(\omega_n \tau_n \gg \omega \tau_s\), we can assume \(\omega \ll \omega_n^2 \tau_n \ll \omega_s^2 \tau_n\) in Eq. (20), and we find that the AV mode is not renormalized in the presence of semiconducting layer (in the first-order perturbation theory with respect to the parameter \(\omega_n^2 \tau_n/\omega\)). One can find the next-order corrections with respect to \(1 \ll \omega_n^2 \tau_n/\omega \ll \omega_s^2 \tau_n/\omega\); however, they are small and we will not discuss them further.

B. Modes interaction if \(v_n < v_s\)

In this case, the density of electrons in the semiconducting layer is smaller than the total density of electrons in the SC layer, \(N_n < N_s\). The intersection of the plasmon mode (of semiconducting layer) with both qAB and AV modes (of SC layer) is possible, depending on the system parameters, as it is shown by dashed curves in Fig. 4 and Fig. 5.

We analyze the interaction between the modes numerically (see solid curves presented in Fig. 4 and Fig. 5). The main conclusion is that the plasmon mode of semiconducting layer interacts with qAB mode stronger than with AV mode since in the latter case, the modes repel each other shifting their positions in the \(q\)-axis. This situation can be explained if we recall that the qAB mode corresponds the oscillation of SC elections and it is accompanied by the oscillations of induced electric fields. In contrast, the AV mode corresponds to the anti-phase oscillations of SC and normal electrons in SC layer resulting in the quasi-neutrality. Therefore, since the interaction between the layers has Coulomb nature, it is not surprising that the plasmon mode interacts with qAB mode much stronger.

IV. ELECTROMAGNETIC FIELD POWER ABSORPTION

Let us expose the system to an external EM field and study light-matter interaction. We will do a numerical analysis of EM power absorption \(\text{Re} (j_n E_0^*)\) [where \(j_n\) is defined in Eq. (4)] in the case when \(v_n < v_s\) and, thus \(N_n < N_s\). In the opposite regime, the interaction of the modes is weak.

First, we switch off the interaction between the layers. The EM power absorption of an isolated semiconducting layer is characterized by a standard Lorenz-shaped resonance at the plasmon frequency. A more interesting situation arises in the isolated SC layer due to the presence of both the qAB and AV modes. The AV mode is quasi-neutral, and hence it interacts with the external EM field weakly, thus the EM power absorption is negligibly small. Indeed, the approximate analytical solution Eq. (12) if substituted into Eq. (9) gives \(\sigma_s = 0\) and \(j_n = 0\) reflecting the quasi-neutral nature of AV mode and nearly zero EM power absorption. The exact numerical calculation of the power absorption in the absence of interaction between SC and semiconducting layers (\(d \to \infty\)), presented in Fig. 5(a), supports this conclusion. At large frequencies \(\omega \gg ku\), the absorption in Fig. 5(a) reaches a plateau with the value equal to the static Drude absorption of EM radiation of normal elec-
trons in SC layer, as it is seen from Eq. (9). Indeed, at $\omega \gg k u$ the conductivity of SC layer resembles the static Drude conductivity $\sigma_s \approx \sigma_{s0}$. The arrow in Fig. 6(a) indicates EM power absorption at the position of AV mode which is negligibly small in comparison with the power absorption corresponding to the plateau value.

In contrast, the qAB mode strongly interacts with external EM field and demonstrates a pronounced peak of absorption with the amplitude much greater than the plateau value, as it is shown in Fig. 6(b).

Furthermore, let us switch the interlayer interaction on and consider the EM power absorption by the hybrid system when the collective modes of SC are hybridised with the plasmon mode in the semiconducting layer (Fig. 7). Figure 7(a) shows the EM power absorption by the hybrid AV-plasmon modes. Weak hybridization of the modes results in two new eigen modes. One of them almost resembles the original plasmon mode, and the other mostly possesses the properties of bare AV mode. From our previous discussion of the light absorption by the bare modes, we expect that the quasi-plasmon mode (the hybrid mode which mostly inherited the properties of the plasmon mode) will give a pronounced resonance in EM absorption spectrum, whereas the other, quasi-AV mode (the hybrid mode which mostly inherited the properties of the AV mode) will not be seen. This reasoning is supported by the exact numerical calculations, shown in Fig. 7(a). We see a single resonance at the frequency of the quasi-plasmon mode with a shifted position with respect to bare plasmon dispersion and a pronounced Lorenz shape. The quasi-AV mode is not seen in the figure.

At large $\Omega$, the EM power absorption spectrum again reaches the plateau with the value equal to the Drude conductivity of the SC layer. The amplitude of the Lorenz-shaped resonance is smaller than the value at the plateau since the total density of electrons in the semiconducting layer is smaller than the density of electrons in SC layer, $N_n < N_s$.

Finally, EM power absorption by the hybrid qAB-
plasmon mode demonstrates a double-resonance structure, indicating that both the new modes formed from the original qAB and plasmon modes (before the coupling) are extremely sensitive to external EM radiation. We want to point out, that such shape of the spectrum can be attributed to the Fano resonance, which is expected in hybrid Bose-Fermi systems [15]. And again, at large frequencies, the energy absorption reaches the plateau with the value equal to the Drude conductivity absorption of SC layer.

V. CONCLUSION

We have studied analytically and numerically the coupling of collective modes in a hybrid two-dimensional electron gas–superconductor structure in the vicinity of the critical temperature of superconducting transition of the superconductor. The layer containing the electron gas in normal state can be represented by a semiconductor or a metallic layer, used in hybrid NS contacts [12]. We have shown, that the superconducting layer lodges two collective modes with drastically different physical properties, which results in their different sensitivity to external electromagnetic fields. As a result, these modes couple differently with the plasmon mode of the normal electron gas spatially separated from the superconducting layer. It manifests itself in different spectra of the electromagnetic power absorption of the hybrid system.

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[1] J. Bardeen, L. N. Cooper, and J. R. Schrieffer, Phys. Rev. 108, 1175 (1957).
[2] P. W. Anderson, Phys. Rev. 112, 1900 (1958).
[3] N. N. Bogoliubov, V. V. Tolmachev, and D. N. Shirkov, New method in the theory of superconductivity (English translation: Consultants Bureau Enterprises, Inc., New York, 1958).
[4] V. G. Vaks, V. M. Galitskii, A. I. Larkin, JETP 14(5), 1177 (1962). [Russian ZhETF, 41(5), 1655 (1962)].
[5] P. I. Arseev, S. O. Loiko, N. K. Fedorov, Phys. Usp. 49, 1 (2006).
[6] R. V. Carlson and A. M. Goldman, Phys. Rev. Lett. 31, 880 (1973).
[7] R. V. Carlson and A. M. Goldman, Phys. Rev. Lett. 34, 11 (1975).
[8] R. V. Carlson and A. M. Goldman, J. Low Temp. Phys. 25, 67 (1976).
[9] A. Schmid and G. Schon, Phys. Rev. Lett. 34, 941 (1975).
[10] S. N. Artemenko and A.F. Volkov, JETP 42 (5), 896 (1975).
[11] S. N. Artemenko and A. F. Volkov, Sov. Phys. Usp. 22, 295 (1979).
[12] B. N. Narozhny, I. L. Aleiner, and B. L. Altshuler, Phys. Rev. B 60, 7213 (1999).
[13] O. Cotlet, S. Zeytinoglu, M. Sigrist, E. Demler, and A. Imamoglu, Phys. Rev. B 93, 054510 (2016).
[14] V. M. Kovalev, A. V. Chaplik, JETP Lett. 98, 331 (2013).
[15] M. V. Boev, V. M. Kovalev, and I. G. Savenko, Phys. Rev. B 94, 241408(R) (2016).
[16] K. H. A. Villegas, V. M. Kovalev, F. V. Kusmartsev, I. G. Savenko, Phys. Rev. B 98, 064502 (2018).
[17] K. H. A. Villegas, F. V. Kusmartsev, Y. Luo, and I. G. Savenko, Phys. Rev. Letters 124, 087701 (2020).
[18] Yu. N. Ovchinnikov, JETP 45(2), 404 (1977).
[19] Y. E. Lozovik and S. M. Apenko, J. Low Temp. Phys. 43, 383 (1981).