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Luminescence of a Cooper Pair

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This Letter theoretically discusses the photon emission spectra of a superconducting \( p-n \) junction. On the basis of the second order perturbation theory for electron-photon interaction, we show that the recombination of a Cooper pair with two \( p \)-type carriers causes enhancement of the luminescence intensity. The calculated results of photon emission spectra explain characteristic features of observed signal in an recent experiment. Our results indicate high functionalities of superconducting light-emitting devices.

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Light-emitting diodes (LEDs) usually fabricated on semiconductors have been an important element of modern technologies. Recent researches seem to focus on producing a well controlled photon and an entangled photon pair [1,2] for realizing quantum information. Superconducting devices have a great advantage in producing entangled quantum states because of its coherent nature [3–5]. Superconducting LEDs [6] have been originally proposed in the context of superradiation. Today, they are a promising source of an entangled photon pair [7]. A recent theoretical study predicts the Josephson radiation in a superconducting \( p-n \) junction [8]. Thus hybrids of superconductor and semiconductor LED undoubtedly have a possibility to provide a key technology in the next generation.

The radiative recombination of Cooper pairs has been observed recently in a InGaAs/InP \( p-n \) junction attaching onto a superconductor Nb [9]. The electroluminescence becomes drastically large at low temperatures below the superconducting transition temperature \( T_c \) of Nb electrode. Surprisingly the degree of enhancement in the luminescence intensity is 1 order of magnitude. Although the experiment has shown clearly effects of superconductivity on the radiative recombination, the mechanism has been an open question. This Letter theoretically addresses this issue. We study the emission spectra of photon in a superconducting \( p-n \) junction by using the second order perturbation expansion for electron-photon interaction. In the second order, we find that a peculiar recombination process to superconductivity enhances the luminescence intensity. In that recombination process, two electrons recombine with two \( p \)-type carriers as a Cooper pair. The theoretical results explain characteristic features of the experimental findings [9]. This Letter not only figures out a mechanism of the large luminescence intensity but also gives a guide for designing highly functional superconducting light-emitting devices.

Let us consider a \( p \)-type semiconductor-superconductor junction under the applied bias voltage \( eV_{sd} \) as shown in Fig. 1(a). The energy is measured from the horizontal line indicated by “0”. The sign of energy in a \( p \)-type semiconductor is chosen to be opposite to that in a superconductor. We assume that a semiconductor and a superconductor are in their local equilibrium which are characterized by the local chemical potential \( \mu_p \) and \( \mu_n \), respectively. The edges of the conduction and valence bands are \( E_c \) and \( E_v \), respectively. In what follows, we use a unit of \( h = k_B = c = 1 \), where \( k_B \) is the Boltzmann constant and \( c \) is

FIG. 1 (color online). Schematic energy diagram of forward biased \( p-n \) junctions. A theoretical model used for calculation is shown in (a). In (c), a realistic junction in experiments is illustrated. Predicted results of the photon spectra are shown in (b) and (d).
the speed of light. The \(p\)-type semiconductor is described by

\[
H_p = \sum_{k,\sigma} \epsilon_p(k) b_{k,\sigma}^\dagger b_{k,\sigma},
\]

where \(\epsilon_p(k) = k^2/(2m_p) + E_v + eV_{sd}/2, m_p\) is the effective mass, and \(b_{k,\sigma}^\dagger(b_{k,\sigma})\) is the creation (annihilation) operator of a \(p\)-type carrier with a wave number \(k\) and spin \(\sigma = \uparrow\) or \(\downarrow\). The photon states are described by

\[
H_{ph} = \sum_q \omega_q \left(a_q^\dagger a_q + \frac{1}{2}\right).
\]

where \(a_q^\dagger(a_q)\) is the creation (annihilation) operator of a photon with a wave number \(q\) and an energy \(\omega_q\). The normal state in a metal is described by

\[
H_{nn} = \sum_{k,\sigma} \left(\frac{k^2}{2m_n} + E_c + eV_{sd}/2\right) c_{k,\sigma}^\dagger c_{k,\sigma},
\]

where \(c_{k,\sigma}^\dagger(c_{k,\sigma})\) is the creation (annihilation) operator of an \(n\)-type carrier and \(m_n\) is the effective mass. The electron-photon interaction Hamiltonian in the dipole approximation is given by

\[
H_I = \sum_{k,q,\sigma} B_{k,q} b_{k-q,\sigma} c_{k,\sigma} a_q^\dagger + \text{H.c.},
\]

where \(B_{k,q}\) is the coupling energy. On the basis of the second order perturbation theory, the number of emitting photon \(N_{ph} = \sum_q a_q^\dagger a_q\) is calculated as

\[
\langle N_{ph} \rangle = \langle N_{ph}(1) \rangle + \langle N_{ph}(2) \rangle,
\]

\[
\langle N_{ph}(1) \rangle = \int_{-\infty}^{t_f} dt_1 \int_{-\infty}^{t_f} dt_2 \langle \chi_0 | H_I^\dagger(t_1) N_{ph} H_I(t_2) | \chi_0 \rangle,
\]

\[
\langle N_{ph}(2) \rangle = \int_{-\infty}^{t_f} dt_1 \int_{-\infty}^{t_f} dt_2 \int_{-\infty}^{t_f} dt_3 \int_{-\infty}^{t_f} dt_4 I(2),
\]

\[
I(2) = \langle \chi_0 | H_I^\dagger(t_1) H_I(t_2) N_{ph} H_I^\dagger(t_3) H_I(t_4) | \chi_0 \rangle,
\]

where \(|\chi_0\rangle \mapsto |0\rangle \otimes |N\rangle \otimes |P\rangle\).

The BCS theory describes superconducting states,

\[
H_{ns} = \sum_{k,\sigma} E_k y_{k,\sigma}^\dagger y_{k,\sigma},
\]

where \(E_k = \sqrt{\varepsilon_k^2(k) + \Delta^2}, \varepsilon_k(k) = k^2/2m_n - \mu_n, \Delta\) is the pair potential, and \(y_{k,\sigma}^\dagger(y_{k,\sigma})\) is the creation (annihilation) operator of a Bogoliubov quasiparticle. We try to consider effects of superconductivity through the Bogoliubov transformation [10]. The description in Eq. (10), however, is valid within a small energy scale near the Fermi level which is at \(\mu_n = E_v + eV_{sd}/2 + \mu_n\) measured from 0.

To apply the BCS theory to the present issue, a rule is necessary to describe the operator in the interaction picture. The transformation connects an electron operator and Bogoliubov operators by

\[
c_{k,\sigma}^\dagger(t) = e^{i\mu(t)} (u_k e^{i\xi_p(k)/2} - s_{\sigma} v_k e^{-i\xi_p(k)/2}) c_{-k,\sigma},
\]

in \((\chi_0| \cdots |\chi_0)\), where \(u_k(v_k) = \left( [1 + (-1)\varepsilon_k(k)/E_k]^{1/2}, s_{\sigma} = 1(-1)\right)\), and \(\sigma\) means the opposite spin to \(\sigma\). The thermal average of operators is carried out in the local equilibrium. In a \(p\)-type semiconductor, for instance, the average of operators is calculated in

\[
\langle N_{ph} \rangle = 2\pi \sum_{k,\sigma} |B_{k,q}|^2 f_{k-q} f_{k} \delta(\omega - E_k) + v_k^2(1 - f_k^q)\delta(\omega + E_k),
\]

where \(\omega = \omega_q - E_g - \mu_n - \mu_p - \xi_p(k - q), E_g = E_v + E_c + eV_{sd}, f_k^q = \left(1 - \tanh(E_k/2T)\right)/2, \) and \(f_k = \left(1 - \tanh(\xi_p(k)/2T)\right)/2\). This result recovers the photon spectra in a normal \(p-n\) junction by tuning \(\Delta \rightarrow 0\), which means that \(E_k \rightarrow -\xi_p(k), u_k \rightarrow 0, v_k \rightarrow 1\) for \(k < k_F\) and \(E_k \rightarrow \xi_p(k), u_k \rightarrow 1, v_k \rightarrow 0\) for \(k > k_F\) with \(k_F\) being the Fermi wave number satisfying \(k_F^2/2m_n = \mu_n\). The threshold of spectra is \(E_g\) and the width of spectra is given by \(\mu_n + \mu_p\). The spectra in Eq. (13) have a broad profile reflecting the quasiparticle density of states as schematically shown in Fig. 1b. Equation (13) depends on temperature through the Fermi distribution function and the pair potential. We have, however, numerically confirmed that the dependence is very weak and that Eq. (13) is almost constant below and near above \(T_c\). The general features of spectra are determined by the energy scales such as \(E_g, \mu_n, \mu_p\). They are much larger than \(\Delta\).

The results of the second order perturbation are given by

\[
I(2) = \sum_{k_1\cdots k_4,q_1\cdots q_4,\sigma_1\cdots \sigma_4} e^{-i\Omega_1 t_1-i\Omega_2 t_2+i\Omega_3 t_3+i\Omega_4 t_4} x B_{k_1,q_1} B_{k_2,q_2} B_{k_3,q_3} \sum_{Q_P Q_P Q_N} Q_{P}\bar{Q}_{Q} \bar{Q}_{N},
\]

with \(\Omega_j(k, q_j) = \omega_{q_j} - \epsilon_F(k - q_j) - \mu_n\). The average of the operators \(\bar{Q}_{P}, \bar{Q}_{Q}, \text{and} \bar{Q}_{N}\) are calculated as follows,

\[
\bar{Q}_{P} = \sum_{q_{5}} \langle 0 | a_{q_1} a_{q_2} a_{q_3} a_{q_4} a_{q_5} | 0 \rangle = 2(\delta_{14} \delta_{23} + \delta_{13} \delta_{24}).
\]
The remaining eight terms in \( \text{QN} \) gives large contribution to the emission spectra at \( \Delta \approx 0 \) denoting the normal density of states in a superconductor. We neglect dependence of \( B \) on wave numbers and assume \( f_{k_f} = 1 \). At \( 1/\tau = 0 \), \( I_0 = \pi N_0 / 2 \Delta \) essentially diverges for small \( \Delta \) with \( N_0 \) denoting the normal density of states in a superconductor.

We first show mathematical reasons of the singularity. Then we will discuss the physics behind the phenomenon. A two-photon emitting process in \( \text{QN}(S) \) is illustrated in Fig. 2(a). The annihilation of a Cooper pair is described by \( c_{-k_1} c_{k_1} \) which includes an operator \( \gamma_{k_1}^\dagger \gamma_{k_1} \). Let us assume that the energy of the initial state is zero. In the first order expansion, the operation of \( \gamma_{k_1}^\dagger \) to the BCS state decreases energy by \( E_k + \mu_n \). At the same time, a \( p \)-type carrier with energy \( E_p(k - q) \) is destructed and a photon with energy \( \omega_k = \epsilon_p(k - q) \) is created. Thus the energy of the intermediate state \( \delta E_1 \) results in \( \delta E_1 = \omega_k - \epsilon_p(k - q) - E_k - \mu_n = \Omega_{k,q} - E_k \). In the perturbation expansion, \( \delta E_1 \) becomes the energy denominator. In the second order, the operation of \( \gamma_{k_1}^\dagger \gamma_{k_1} \) the destruction of a \( p \)-type carrier, and the creation of a photon gain energy by \( E_k - \mu_n \) with energy \( \omega_k - \epsilon_p(k + q) \), respectively. Therefore the difference in energy between the intermediate state and the final one becomes \( \delta E_2 = \omega_k - \epsilon_p(k - q) - E_k - \mu_n = \Omega_{k,q} + E_k \). The perturbation theory requires the energy conservation between the initial state and the final one (i.e., \( \delta E_1 + \delta E_2 = 0 \)), which leads to \( 2 \Omega_{k,q} = 0 \). As a result, a small value of \( E_k \) remains in the denominator as shown in Eq. (20). The physics behind the phenomena is simple. The BCS state has an ability to emit a pair of photons with remaining in its state almost unchanged because the BCS state is the eigenstate of \( \gamma_{k_1}^\dagger \gamma_{k_1} \). The equation \( \Omega_{k,q} = 0 \) describes the emitting condition of two photons. The threshold and width of spectra are \( E_q + \mu_n \) and \( \mu_p \), respectively. In Fig. 1(b), we schematically show predicted spectra in the second order process.

The singular behavior in perturbation expansion implies an importance of higher order terms for predicting the luminescence intensity quantitatively. Here we do not discuss this issue, but choose an alternative way of regularizing the obtained results. In what follows, we introduce a finite relaxation time. First we consider a mean free time..
whose thickness $n$ superconductor is attached to an $T_c$.

In (b), we consider the relaxation due to inelastic scatterings.

The luminescence intensity at $I$ increases $P$ mainly from a quantum well which is sandwiched by the $Cie$.

The amplitude at $I_0$ is a constant of the order of unity. In Fig. 3(a), we consider relax-

de a quantum well can be proportional to $\frac{\Delta^2 e^{-2L_w/\xi_\tau}}{(\omega_q - \omega_0)^2 + (\tau_\tau)^2}$. 

Finally we modify Eq. (22) to describe the photon spectra in realistic junctions as shown in Fig. 1(c). A superconductor is attached to an $n$-type semiconductor whose thickness $L_w$ is about 30–50 nm [9]. The proximity effect enables the penetration of Cooper pairs into the $n$-type semiconductor. In experiments, photons are emitted mainly from a quantum well which is sandwiched by the $p$- and $n$-type semiconductor. The pair amplitude in the quantum well can be proportional to $\Delta e^{-2L_w/\xi_\tau}$ with $\xi_\tau = \sqrt{D}/2\pi T$ and $D$ being the diffusion constant in the $n$-type semiconductor. The quantum well would be replaced by a quantum dot near future. The level in the quantum well (dot) $E_w$ should coincide with the Fermi level in the $n$-type semiconductor $\mu_n$. Namely, $|E_w - \mu_n|$ must be less than both the Thouless energy $E_{\text{Tb}} = D/L_w^2$ and $\Delta$. This resonant condition is particularly important for a Cooper pair to penetrate into the quantum well (dot). The emission spectra have a peak at $\omega_q$ and the peak width is given by $\Gamma = \tau_\tau N_0$, where $\tau_\tau$ is the transfer integral between the quantum well (dot) and the semiconductor. The argument above is summarized by an equation for $T \approx T_c$.

$\frac{I_0}{I_0(0)} = \frac{\alpha + 1}{\alpha - 1} \left( 1 + \frac{\Delta^2}{\xi_\tau^2} \right)$

where we introduce the Lorentz resonant function by hand.

In the experiment, $\xi_\tau$ estimated to be 680 nm at 4 K is much larger than $L_w$ below $T_c$. Thus the theoretical results in Fig. 3 may describe experimental results of the luminescence intensity. In fact, the experimental results of Fig. 6(b) in Ref. [9] show a very similar line shape to that in Fig. 3(a) with $\alpha = 1$.

In conclusion, we have studied the photon emission spectra in a superconducting $p$-$n$ junction based on the second order perturbation theory for electron-photon interaction. We have found in the second order expansion that a peculiar recombination process to superconductivity enhances the luminescence intensity. The theoretical results explain temperature dependence of the luminescence intensity observed in an recent experiment.


[11] The speed of light is much larger than the Fermi velocity. In such case, this condition becomes $|q_1| \sim |q_2|$.