Capacitively coupled hot-electron nanobolometer as far-infrared photon counter

Dragoş-Victor Anghel\textsuperscript{a)}

University of Oslo, Department of Physics, P.O.Box 1048 Blindern, N-0316 Oslo, Norway and NIPNE – “HH”, P.O.Box MG-6, R.O.-76900 Bucureşti - Măgurele, Romania

Leonid Kuzmin\textsuperscript{b)}

Department of Microelectronics and Nanoscience, Physics and Engineering Physics, Chalmers University of Technology, S-41296 Göteborg, Sweden

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We show theoretically that hot-electron nanobolometers consisting of a small piece of normal metal, capacitively coupled to a superconducting antenna through a pair of normal metal–insulator–superconductor (NIS) tunnel junctions may be used as far-infrared photon counters. To make the device most effective at high counting rates, we suggest the use of the bolometer in the simplest configuration, when the NIS tunnel junctions are used as both an electron cooler and thermometer. The absorption of the photon in the normal metal produces a pulse in the electron temperature, which is measured by the NIS junctions. The counter may resolve photons up to 0.3–0.4 mm wavelength and has a typical re-equilibration time constant of about 20 ns.

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The far-infrared region is one of the richest areas of spectroscopic research, with applications ranging from molecular physics to astronomy. These applications, especially the astrophysical observations, require extremely sensitive detectors. Depending on the energy of the incident photons, the input power, and the response time of the detector, the device may resolve individual quanta (quantum detectors), or may measure the total input power (integrating detectors). In general, integrating detectors measure long-wavelength radiation fluxes (infrared and far infrared regions \cite{1,2}), while high energy photons (for example, x ray) are observed by quantum detectors. The first x-ray detector using a normal metal film as an absorber and a normal metal–insulator–superconductor (NIS) tunnel junction as a thermometer, was built by Nahum et al. \cite{3}. The photon interacts with the detector by producing a hot spot in the normal metal. The energy dissipated into the hot spot diffuses then into the whole absorber and the increase of electron temperature is measured by the NIS thermometer. The re-equilibration time constant of the detector was of the order of 10 $\mu$s and determined mostly by the electron–phonon coupling. The same phenomenon of hot spot formation due to the photon absorption was suggested for the detection of longer wavelength photons \cite{4}. Recently, Semenov et al. \cite{5} constructed a near-infrared photon counter based on a current biased narrow superconducting strip. The hot spot formed in the strip leading to a break of the superconductivity across it, which further produced a signal in the voltage read out. At the experimental stage reported until now, the time constant of the detector is of the order of 100 ps and is limited by the amplifier bandwidth—the estimated thermalization time of the superconducting strip being of the order of tens of picoseconds.

To increase the limit wavelength of detectable individual IR photons, at a counting rate high enough for astronomical observations but still within the range of broadband amplifiers, we suggest the use of the capacitively coupled hot-electron microbolometer \cite{6} as a photon counter. The detector (Fig. 1) is formed of a small normal metal island connected to the superconducting antenna by two symmetric NIS tunnel junctions. As in the case of x-ray detectors \cite{3}, the energy of the quanta, dissipated in the normal metal island [which we call the thermal sensing element (TSE)] leads to an increase of the electron temperature, which is measured by the NIS tunnel junctions. The time scale of the heat diffusion into the whole normal metal element is $\tau_d \sim L^2/\left(\pi^2 D\right)$, where $L$ is the linear dimension of the sample and $D$ is the electron diffusion constant. If we set $L = 1 \mu$m, then, for a typical value of $D = 10^{-4} \text{m}^2\text{s}^{-1}$ in Cu thin films,

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure1.png}
\caption{Schematic drawing of the main part of the photon quantum detector.}
\end{figure}

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\textsuperscript{a)}E-mail: dragos@fys.uio.no
\textsuperscript{b)}E-mail: kuzmin@fy.chalmers.se
\( \tau_1 \approx 1 \) ns. To have an accurate reading of the temperature variation in the TSE, the detector re-equilibration time, \( \tau_c \), should be larger than \( \tau_1 \). Although not shown in Fig. \ref{fig:fig1}, the detector (TSE plus antenna) is supported by an insulating substrate, which acts also as a heat bath of temperature \( T_b \). Here, we assume perfect thermal contact between the substrate and the detector lattice, and we set the lattice temperatures of both the TSE and the antenna equal to \( T_b \). On the other hand, in the operating temperature range of the detector (of the order of 100 mK) the electron–phonon coupling is weak and we assign to the electron gas in the TSE an effective temperature \( T_e \), which, in general, is different from \( T_b \). For the most effective operation of the nanobolometer, the NIS junctions are used both to cool the electron gas in the detector below the heat bath temperature \( T_b \) and to measure the temperature pulse (the principles of NIS cooling and temperature measurement are presented in Refs. \ref{ref1} \ref{ref2} and \ref{ref3} respectively). In the normal metal, the power flow from the electron system to the lattice is \( \dot{Q}_{\text{ep}} = \Sigma_{\text{ep}} \Omega(T_e^5 - T_b^5) \), where \( \Omega \) is the volume of the TSE. We approximate the electron specific heat with \( C_V = \Omega(\frac{k_B T_e^3}{3h^3}) \sqrt{2m_e^2} \), where \( \epsilon_F \) is the Fermi energy and \( m_e \) is the electron mass. For the concrete calculations, we shall consider Cu as the normal metal and Al as the superconductor since these two metals are routinely used in the construction of microrefrigerators. The cooling power of one junction is denoted by \( \dot{Q}_J \) and, for a symmetric setup as the one shown in Fig. \ref{fig:fig1}, the total power extracted through the junctions is \( 2\dot{Q}_J \). \( \dot{Q}_J \) is a function of both the effective temperature of the quasiparticles in the superconductor, \( T_s \), and \( T_e \). We assume that the superconductor is thermalized by using efficient traps \ref{ref4}, so we use also \( T_s = T_b \). At equilibrium (no incident radiation), the total power extracted from the normal metal is zero:

\[
\dot{Q}_T = 2\dot{Q}_J + \dot{Q}_{\text{ep}} = 0. \tag{1}
\]

Equation (1) sets the equilibrium electron temperature, \( T_{e0} \), as a function of \( T_b \). As a thermometer, the NIS junctions may work in either voltage biased (VB) or current biased (CB) regime. For a linear response of the detector, the temperature pulse due to the photon absorption, \( \delta T_e \equiv T_e - T_{e0} \ll T_e \), has an exponential decay \( \delta T_e(t) = \delta T_e(0) \exp(-t/\tau) \), where \( \delta T_e(0) = \hbar \omega/C_V \) and we assume that \( \tau \equiv C_V \left( \frac{\partial \dot{Q}_T}{\partial T_e} \right)^{-1}_{T_s = T_{e0}} \gg \tau_1 \). Using Eq. (1) and the expression for \( \tau \), we write \( \tau^{-1} = \tau_1^{-1} + \tau_{\text{ep}}^{-1} \) in obvious notations. For Cu at \( T_b = 100 \) mK (\( \Sigma_{\text{ep}} = 4 \) nWK\(^{-5}\)\(\mu m^{-3} \)), \( \tau_{\text{ep}} \approx 3.5 \times 10^{-5} \) s, which will turn out to be much larger than \( \tau_1 \).

Besides \( \tau \), the most important figure of merit of the detector is the energy resolution (ER). A photon can be detected if the signal produced by its absorption is larger than the square root of the mean square fluctuation of the measured quantity, \( \langle \delta^2 M \rangle^{1/2} \equiv \langle (M - \langle M \rangle)^2 \rangle^{1/2} \), where \( M \) is current (VB) or voltage (CB), and by \( \langle \cdot \rangle \) we denote the statistical average. The ER can then be defined as \( \langle \delta^2 M \rangle^{1/2}/m_{\text{max}} \), where \( m_{\text{max}} \) is the amplitude of the measured pulse. Due to the finite bandwidth of the measuring device, \( \tau_1^{-1} \), the measured and real values of \( M \), say \( m_{\text{meas}}(t) \) and \( m(t) \), respectively, may be related in a general formalism by the equation \( dM_m(t) = \tau_1^{-1} [M_m(t) - M(t)] \, dt \). This leads to the usual frequency dependent amplification factor, proportional to \( [1 + (\omega \tau_c)^2]^{-1/2} \). Therefore, the measured mean square fluctuation is \( \langle \delta^2 M_m \rangle = \int_0^\infty \langle \delta^2 M_c \rangle [1 + (\omega \tau_c)^2] d\omega \), where \( \langle \delta^2 M_c \rangle \) is the spectral density of noise. In VB and CB regimes, \( M_c \equiv I_1 \) (current through the junctions) and \( M_c = 2V \) (voltage across both junctions), respectively. Following Ref. \ref{ref5} and disregarding the external circuit of the detector, we write the total fluctuation \( \delta I_J \) (VB) as the superposition of the fluctuation due to the discrete transport of charges through the junctions, and the current fluctuations induced by the temperature and particle fluctuations in the normal metal island. In the CB regime, the voltage fluctuation is induced only by the particle number fluctuation. Using this, we obtain the following spectral densities:

\[
\langle \delta^2 I_J \rangle_\omega = \langle \delta^2 I_{J,\text{shot}} \rangle_\omega + \left( \frac{\partial I_J}{\partial T_e} \right)^2 \langle \delta^2 T_e \rangle_\omega
\]

\[
+ \frac{1}{\epsilon^2 F^2} \left( \frac{\partial \epsilon_F}{\partial N} \right)^2 \langle \delta^2 T_e \rangle_\omega
\]

\[
+ 2 \frac{\partial I_J}{\partial \epsilon} \frac{\partial \epsilon}{\partial T_e} Re \left( \frac{\delta T_e}{\delta N} \right) \langle \delta^2 I_{J,\text{shot}} \rangle_\omega.
\]

\[
\langle \delta^2 V \rangle_\omega = \frac{1}{\epsilon^2 F^2} \left( \frac{\partial \epsilon_F}{\partial N} \right)^2 \langle \delta^2 I_{J,\text{shot}} \rangle_\omega,
\]

where \( \epsilon_F \) is the Fermi energy, \( N \) is the number of electrons in the normal metal, and \( \langle \delta^2 N \rangle_\omega = \langle \delta^2 I_J \rangle_\omega \cdot (\omega)^{-2} \). We also assumed that \( \partial (\epsilon F)/\partial N \equiv \epsilon_F/\partial N = (2/3) (\epsilon_F/N) \), for the ideal gas. The method to evaluate \( \langle \delta^2 T_e \rangle_\omega \) and the correlation terms was given in Ref. \ref{ref6}. In each case, the amplifier noise should be added quadratically to the fluctuations above. \( \langle \delta^2 I_{J,\text{shot}} \rangle_\omega \) is the current Poissonian shot noise (white noise). According to Eqs. \ref{eq:1} and \ref{eq:2}, in both VB and CB regimes \( \langle \delta^2 M_c \rangle \propto \omega^{-2} \) for \( \omega \to 0 \), which implies an “infrared” divergence of the total measured fluctuation \( \langle \delta^2 M_m \rangle \). Apparently, this fact would make our approach totally hopeless. Fortunately, we can avoid this, at least for the VB regime. From Eq. \ref{eq:1} we notice that \( \langle \delta^2 I_J \rangle_\omega \) is the sum of an \( \omega \)-independent term, call it \( \langle \delta^2 I_{\text{J,0}} \rangle_\omega \), and an \( \omega \)-dependent one, say \( \langle \delta^2 I_{\text{J,1}} \rangle_\omega \). Obviously, \( \langle \delta^2 I_{\text{J,1}} \rangle_\omega \to 0 \) or \( \omega \to 0 \) or \( \omega \to 0 \), respectively. We can define the crossover between \( \omega \)-independent and \( \omega \)-dependent regimes as \( \langle \delta^2 I_J \rangle_\omega \propto \langle \delta^2 F \rangle_\omega \propto \langle \delta^2 I_J \rangle_\omega \propto \omega^{-2} \). For \( \omega \ll 1/\tau \) and since “slow” fluctuations of \( I_J \) do not influence the reading of the signal due to the photon absorption, we can approximate the total fluctuation by using a convenient
cutoff of the integral at the lower end:

$$\langle \delta^2 I_{\text{m}} \rangle \approx \int_{\omega_0}^{\infty} \frac{\langle \delta^2 I_\omega \rangle}{1 + (\omega/\omega_c)^2} \, d\omega \leq \frac{\pi}{2} \frac{\langle \delta^2 I_0 \rangle}{\tau_c},$$

(4)

where $\omega_c \ll \omega_0 \ll 1/\tau$. From Eq. (3), we see that this procedure cannot be applied to CB regime. Moreover, evaluations of $\langle \delta^2 I_{\text{m}} \rangle$ and $\langle \delta^2 V_{\text{m}} \rangle$, using a similar cutoff show that the second regime is not practical, and in what follows we shall consider only the VB regime.

Let us focus now on the signal measurement. In the linear regime, if the photon is absorbed at $t = 0$, then $I_{\text{f}}(t < 0) \equiv I_{\text{f}}(T_{\text{co}})$, while $I_{\text{f}}(t \geq 0) = I_{\text{f}}(T_{\text{co}}) + i_0 e^{-t/\tau}$, where $i_0 = I_{\text{f}}[T_{\text{e}}(0)] - I_{\text{f}}(T_{\text{co}})$. If $\tau_c \gg \tau$, the measured current is $I_{\text{m}}(t) = I_{\text{f}}(T_{\text{co}}) + i_m(t)$, where $i_m(t < 0) = 0$ and $i_m(t \geq 0) = i_0 \tau_c (\tau_c - \tau)^{-1} e^{-t/\tau} - e^{-t/\tau}$. If we denote $x \equiv \tau_c / \tau$, then the maximum value of $i_m(t)$ is $i_{\text{max}} = i_0 x^{1/2} (1 - x)$ and provides a way to calculate the energy of the absorbed photon. According to the definition, the relative error in the determination of $i_0$, $\lambda = (\delta^2 I_{\text{m}})^{1/2}/i_{\text{max}}$ is also the ER. If $h\omega_{\text{ph}}$ is the energy of the incoming photon, in the linear approximation $i_0 \approx (\partial I_{\text{f}}/\partial T_{\text{e}})_{T_{\text{co}}} [h\omega_{\text{ph}}/C_{\text{V}}(T_{\text{co}})] \ll I_{\text{f}}(T_{\text{co}})$. Plugging this together with (4) in the expression for ER, we get

$$\text{ER} = x^{1/2} \sqrt{\frac{\pi\langle \delta^2 I_0 \rangle}{2\tau_c} \frac{C_{\text{V}}(T_{\text{co}})}{\langle \partial I_{\text{f}}/\partial T_{\text{e}} \rangle_{T_{\text{co}}} h\omega}}.$$  

Since $x^{-(1+x)/2(1-x)}$ has a minimum equal to $e$ for $x = 1$, we obtain the optimum energy resolution, $\text{ER}_{\text{opt}} = e \left[ (\pi/2)\langle \delta^2 I_0 \rangle / \tau_c \right]^{1/2} C_{\text{V}}(T_{\text{co}}) \left[ \langle \partial I_{\text{f}}/\partial T_{\text{e}} \rangle_{T_{\text{co}}} h\omega \right]^{-1}$.

We now turn to concrete calculations. For this, we take a typical value: $T_{\text{b}} = 300 \text{ mK}$, and we set our goal to $T_{\text{co}} = 100 \text{ mK}$. If we set $\Omega = 2 \times 10^{-3} \text{ m}\mu^3$, then $T_{\text{e}}$ increases to about 200 mK after the absorption of a photon of 100 $\mu$m wavelength. With all these parameters fixed, Eq. (3) is an equation in the junction tunnel resistance, $R_T$, and $V$. Using this equation, we write $\text{ER}_{\text{opt}}$ as a function of $V$. This function has a minimum at $V = V_{\text{opt}} \approx 0.85 \Delta / e$ ($\Delta$ is the gap energy in the superconductor). Setting the bias voltage at $V_{\text{opt}}$, we obtain $R_T \approx 5.4 \text{ k}\Omega$ (Eq. (3)) and $\tau \approx 17 \text{ ns}$. Moreover, in this case $\tau_c \approx 1/23 \ll 1$ which justifies the assumption behind Eq. (3) and, therefore, the expression for $\text{ER}_{\text{opt}}$. In Fig. 2, we show $\text{ER}_{\text{opt}}(V)$, which was calculated using $\langle \delta^2 I_{\omega < \tau_c} \rangle$ instead of $\langle \delta^2 I_0 \rangle$ to set an upper limit for nonlinear effects in $\text{ER}_{\text{opt}}$ as a function of the incident photon wavelength. From this plot, it appears that the device may work as a counter up to about 0.4 mm photon wavelength. For longer wavelengths, $\text{ER}_{\text{opt}}$ is above 1, so, in principle, the signal may not be distinguished from the noise. Nevertheless, the performances may be improved by decreasing $T_{\text{b}}$ and $T_{\text{co}}$. If $\delta T_{\text{e}} \ll T_{\text{co}}$ is not satisfied, the response time would show a variation with the temperature. If we define the temperature dependent re-equilibration time, $\tau'(T_{\text{co}}) \equiv -(T_{\text{co}} - T_{\text{e}})/(dT_{\text{e}}/dt)$, we observe that this does not vary much in the range of interest (see inset of Fig. 2). The approximation

$$\tau'(T_{\text{co}}) \approx \tau(T_{\text{co}}) = \tau$$

was used for the convenience of performing analytical calculations. For a concrete experimental set-up one can calculate numerically the energy resolution, taking into account nonlinear effects. According to the evaluations above, the results should not differ significantly.

In conclusion, we presented a far-infrared photon counter formed of a small piece of normal metal, capacitively coupled to a superconducting antenna via two NIS tunnel junctions. The photon energy is released in the normal metal and the junctions serve as both an electron cooler and thermometer. At a bath temperature $T_{\text{b}} = 300 \text{ mK}$, the electron temperature in the normal metal may be reduced to $T_{\text{co}} = 100 \text{ mK}$ by the cooling effect of the junctions, which tunes the response time of the detector (time constant for the detector re-equilibration after one photon absorption) to $\tau \approx 17 \text{ ns}$ and the tunnel resistance of each junction to $R_T \approx 5.4 \text{ k}\Omega$. For a volume of the normal metal $\Omega = 2 \times 10^{-3} \text{ m}\mu^3$, the calculated ER is 0.23 at a photon wavelength $\lambda = 0.1 \text{ mm}$ (Fig. 2). The ER reaches 1 (the counting limit) for $\lambda \approx 0.4 \text{ mm}$. To allow for counting longer-wavelength photons, one should retune the detector parameters. The technological and physical difficulties raised eventually by the construction of small tunnel junctions with such low transparencies may be overcome either by the use of ferromagnetic materials to suppress the Andreev current across the too thin barriers, as suggested by Giazotto et al. [13], or by the use of heavily doped semiconductors, which would allow one to increase the size of the absorber [14].

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