

Graphene-insulator-superconductor junctions as thermoelectric bolometers

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We design a superconducting thermoelectric bolometer made of a graphene-insulator-superconductor tunnel junction. Our detector has the advantage of being passive, as it directly transduces input power to a voltage without the need to modulate an external bias. We characterize the device via numerical simulation of the full nonlinear thermal dynamical model of the junction, considering heating of both sides of the junction. While estimating noise contributions, we found expressions due to the temperatures of both sides being different from the bath temperature. Numerical simulations show a noise equivalent power of approximately $4 \times 10^{-17} \text{ W}/\sqrt{\text{Hz}}$ for an input power of approximately 10^{-16} W , a response time of approximately 200 ns, and an integration time to obtain a signal-to-noise ratio of 1 of approximately 100 μs for an input power of approximately 10^{-13} W . Therefore, the device shows promise for large-array cosmological experiment applications, also considering its advantages for fabrication and heat budget.

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I. INTRODUCTION

Research on terahertz detection has grown rapidly in recent years due to its relevance in both fundamental science [1,2] and a wide range of technological applications [3–5]. In this frequency range, thermal detectors such as bolometers become especially important because of the so-called terahertz gap—namely, the lack of reliable solid-state devices based on direct photon absorption [6]. For cosmic microwave background measurements, present requirements for noise equivalent power (NEP) and saturation power P_{sat} typically lie around approximately $10^{-17} \text{ W}/\sqrt{\text{Hz}}$ and approximately 10 pW, respectively [7–9], whereas upcoming experiments aim to introduce stricter requirements, with $\text{NEP} \sim 3 \times 10^{-20} \text{ W}/\sqrt{\text{Hz}}$ and $P_{\text{sat}} \sim 0.2 \text{ fW}$ [10]. The transition-edge sensor (TES) remains the benchmark technology for bolometry in astroparticle physics [9,11], with a measured best performance of $\text{NEP} = 3 \times 10^{-19} \text{ W}/\sqrt{\text{Hz}}$ [12]. However, its operation is complicated by the difficulty of dc readout, by the need for biasing electronics, and by the heating associated with it [13], while for ground-based experiments the NEP is still limited to approximately $10^{-17} \text{ W}/\sqrt{\text{Hz}}$ by environmental factors [9].

An alternative approach is offered by kinetic inductance detectors (KIDs), showing a record NEP of approximately $3 \times 10^{-19} \text{ W}/\sqrt{\text{Hz}}$ [14]. They have already been deployed—for example, in the OLIMPO experiment, achieving $\text{NEP} = 4.5 \times 10^{-17} \text{ W}/\sqrt{\text{Hz}}$ at 150 GHz with an input power of 3 pW [15]. KIDs alleviate the readout challenges and mitigate heating issues typical of TES-based systems. However, their physics is quite complex, and they require costly electronics and wiring for probe signals. They can absorb photons directly by destroying a Cooper pair, reducing the need for antennas, but only above the superconducting gap $\nu > \Delta$ [16] (for Al, $\nu > 90 \text{ GHz}$), which poses limitations for CMB applications.

Another class of detectors uses superconducting tunnel junctions (STJs). Their working principle relies on the strong temperature dependence of STJ conductivity, and various junction types have been explored, including metal-insulator-superconductor devices used in hot-electron bolometers [17] and cold-electron bolometers [18,19]. Both these designs require voltage or current biasing, with the same problems as mentioned for the other technologies.

A different approach to detection is represented by passive detectors, which directly generate the signal from the input power without the need to modulate a bias. A passive detector concept is the superconducting thermoelectric bolometer [13], which measures the thermoelectric current or voltage produced by heating induced by the incident signal. Such a detector could, in principle, overcome several of the issues described above. Among the possible implementations,

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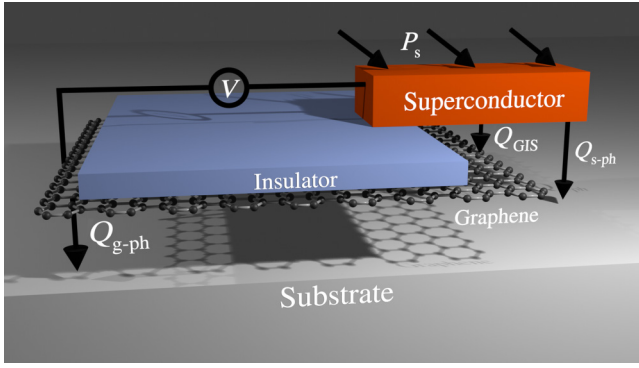


FIG. 1. Depiction of the thermoelectric GIS bolometer schematizing its thermal dynamics and the measurement setup.

STJ-based structures are particularly promising, as strong thermoelectric effects arise in ferromagnet-insulator-superconductor (FIS) junctions [20]. A further mechanism for strong thermoelectricity, accessible in superconductor-insulator-superconductor' (SIS') junctions [21], is nonlinear thermoelectricity, obtained by heating the superconductor with the larger Δ . Peculiarly, the induced thermovoltage scales with the superconducting gap Δ rather than with the temperature T_h of the hotter electrode [21]. This enables sizable thermovoltages even for weak heating, since a small increase in T_h can still drive a signal $\sim \Delta$. The main limitation of FIS and SIS' junctions, however, is the requirement of a strong magnetic field within the junction—needed either to activate thermoelectricity (FIS) or to suppress the Josephson current (SIS').

In this article, we propose the design of a thermoelectric bolometer made of a graphene-insulator-superconductor (GIS) tunnel junction (Fig. 1). There are already several examples of GIS junction-based bolometers in the literature [22–24], but they generate signals from the temperature dependence of resistance measured via biasing and not from thermoelectricity generation.

We begin by describing the thermal dynamics of the device by using a numerical nonlinear model for thermoelectricity in superconducting junctions introduced in Ref. [25], and then we discuss the noise sources we considered and how we estimated their contributions to the total noise of the device. Finally, we present the main features and figures of merit of the device—namely, the responsivity, the noise equivalent power, and the characteristic times.

II. SIGNAL MODEL AND RESULTS

When one side of a GIS junction heats up, a thermoelectric current starts to flow through it [26]. Therefore, we can design a thermal radiation detector by coupling an antenna [27] or an absorber [28] to one side of the junction and measuring the generated closed-circuit Peltier current

I_0 or the open-circuit Seebeck potential V_S . The thermoelectric current or potential generated by the heat carried by the impinging radiation represents the entire signal of our device. Intuitively, this means that we need the largest possible difference between the superconductor temperature T_s and the graphene temperature T_g , thus implying that we will need at least one temperature to be much greater than the bath temperature T_b . As shown in Ref. [25], linear transport models fail to describe the thermal behavior of the system in this regime. Following that discussion, we use the nonlinear numerical model defined in Ref. [25] to describe the junction behavior and estimate the noise performance of the detector. Here we summarize only its main points.

Current and heat transport across the GIS junction can be described by the usual tunneling formulas [25,29]

$$I(V, T_s, T_g) = \frac{1}{eR_T} \int dE \rho_s(E, T_s) \times \rho_g(E - E_F - eV) [f(E - eV, T_g) - f(E, T_s)], \quad (1)$$

$$Q_{GIS}(V, T_s, T_g) = \frac{1}{e^2 R_T} \int dE (E - eV) \rho_g(E - E_F - eV) \times \rho_s(E, T_s) [f(E, T_s) - f(E - eV, T_g)], \quad (2)$$

where R_T is the tunneling resistance, $f(E, T)$ is the Fermi function, V is the electrostatic potential across the junction, E_F is the Fermi level, and ρ_s and ρ_g are the density of states (DOS) of the superconductor and the DOS of graphene, respectively, defined by Eqs. (2c), (4), and (8) in Ref. [25]. We consider here doped graphene with electron density $n_0 = 1 \times 10^{12} \text{ cm}^{-2}$, corresponding to $E_F = 98 \text{ meV}$. The effects of different values of E_F were explored in Ref. [26]. We also choose Al as the superconductor, with DOS at the Fermi level in the metallic state $N_F = 2.15 \times 10^{47} \text{ J}^{-1} \text{ m}^{-3}$ [30], and zero-temperature energy gap $\Delta_0 = 2 \times 10^{-4} \text{ eV}$.

From these expressions, we can obtain $I_0 = I(0, T_s, T_g)$ and V_S by solving $I(V_S, T_s, T_g) = 0$ for any T_s and T_g . In both open-circuit and closed-circuit operation, we assume that we have an ideal voltmeter and an ideal ammeter connecting the superconductor and the graphene layers. Therefore, to compute the response of our device to an input power, we need to describe the thermal dynamics of the junction. Each side of the junction exchanges a power Q_{GIS} with the other side, following Eq. (2), and dissipates power $Q_{s-ph}(T_s)$ [Eq. (6) in Ref. [25]] and $Q_{g-ph}(T_g)$ [Eq. (7a) in Ref. [25]] through interaction with the phonon gas, which is at the bath temperature T_b if we assume a negligible Kapitza resistance [24]. The physics translates into a system of two coupled nonlinear differential equations for the time-dependent temperatures $T_s(t)$, $T_g(t)$ that we solve

numerically [25]:

$$\begin{cases} C_s(T_s) \frac{dT_s}{dt} = -Q_{\text{GIS}}(T_s, T_g) - Q_{\text{s-ph}}(T_s) + P_s, \\ C_g(T_g) \frac{dT_g}{dt} = Q_{\text{GIS}}(T_s, T_g) - Q_{\text{g-ph}}(T_g) + P_g, \end{cases} \quad (3)$$

where $C_g(T) = A_g(\pi^2/3)k_B^2 n_g(E_F)T$ [24] is the thermal capacitance of the graphene layer, with A_g the area of the graphene layer, $n_g(E_F)$ the density of electrons at the Fermi level, and k_B the Boltzmann constant, and $C_s(T) = TdS_s/dT$ is the thermal capacitance of the superconductor layer, computed from the quasiparticle entropy $S_s(T)$ [31]. P_s and P_g are the external radiant powers impinging on the superconductor and the graphene. We explore only results with $P_g = 0$ and $P_s \neq 0$, because with $T_g > T_s$ the absence of nonlinear thermoelectricity reduces the signal [25], as also explained below. We also restrict our analysis to the open-circuit configuration because nonlinear thermoelectricity in this system induces relatively larger V_S than I_0 [25]. We choose a specific geometry with $A_g = 100 \mu\text{m}^2$, a superconductor volume $\mathcal{V}_s$ of $0.5 \times 10^{-3} \mu\text{m}^3$, and overlap area $A_j = 0.01 \mu\text{m}^2$ because this geometry allows the formation of the largest thermal gradient $\delta T = T_s - T_g$ and gives the best results in terms of signal and noise. We use this geometry from now on. Some effects of different geometries are analyzed in Ref. [25].

For a bolometer, we can assume that the variation of P_s is slower than the system response time τ_{th} (and we later prove that τ_{th} is reasonably small enough), allowing us to use the steady-state values $T_s(t \rightarrow \infty) \equiv T_s$ and $T_g(t \rightarrow \infty) \equiv T_g$ to compute all physical quantities. We numerically solve Eq. (3) for different T_b values to obtain the dependence of T_s and T_g on P_s , shown in Fig. 2(a). As can be seen from Fig. 2(a), δT is monotonic in P_s , but the dependence is not linear as predicted by linear models [13], even for very small P_s [25]. We can also see the effect of the difference in T_b vanishing for larger P_s , where dissipation processes dominate the thermal dynamics of the system.

The most representative figure of merit for the signal is the responsivity $\mathcal{R} \equiv \delta V_S / \delta P_s$, defined as the variation in the signal (in open-circuit δV_S) induced by a small variation in the input power δP_s . Since we consider this small variation to happen over a background power P_s , we can compute the system behavior by linearizing Eq. (3) around the steady-state solutions $\{T_s(P_s), T_g(P_s)\}$. After a Fourier transform $t \rightarrow \omega$, we can describe the system's linear response $\{\delta T_s, \delta T_g\}$ around a point $\{T_s, T_g\}$ to power

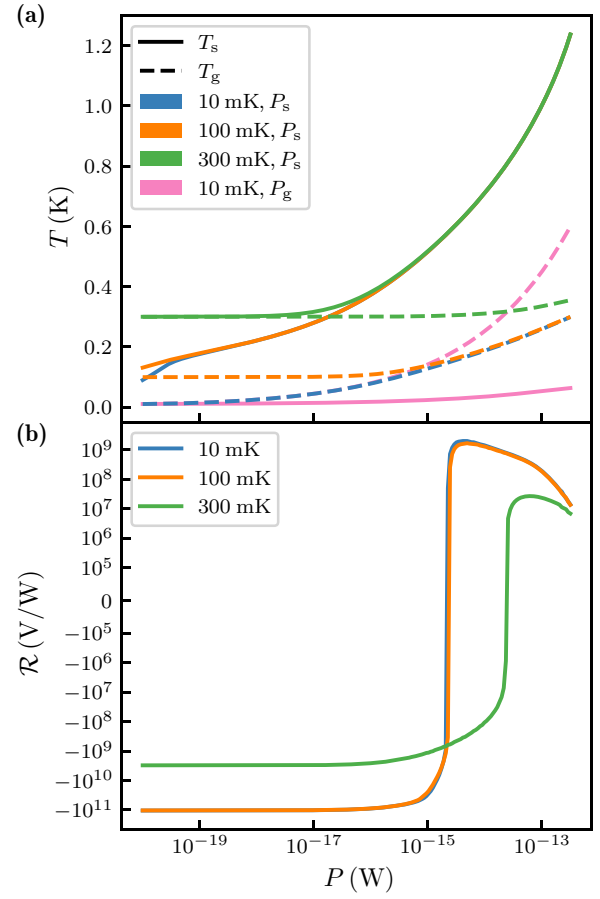


FIG. 2. Device response to steady input power P . (a) Dependence on P of superconductor and graphene temperatures T_s and T_g for different T_b values and signal application sides. When the superconductor is heated, $T_s(P_s)$ converges to the same curve for all T_b values at large values of P_s . Application of the signal to graphene generates a smaller thermal gradient. (b) Responsivity $\mathcal{R}$ dependence on P for different T_b values with power impinging on the superconductor. $\mathcal{R}$ shows large negative values for small P , then becomes positive, passing through zero. This is due to a saturation effect described in the main text and in Ref. [25].

oscillations $\delta P_s(\omega)$ and $\delta P_g(\omega)$ by linearizing Eq. (3) as

$$\begin{cases} i\omega C_s \delta T_s = -\left(\frac{\partial Q_{\text{GIS}}}{\partial T_s} + \frac{\partial Q_{\text{s-ph}}}{\partial T_s}\right) \delta T_s - \frac{\partial Q_{\text{GIS}}}{\partial T_g} \delta T_g + \delta P_s, \\ i\omega C_g \delta T_g = \left(\frac{\partial Q_{\text{GIS}}}{\partial T_g} - \frac{\partial Q_{\text{g-ph}}}{\partial T_g}\right) \delta T_g + \frac{\partial Q_{\text{GIS}}}{\partial T_s} \delta T_s + \delta P_g. \end{cases} \quad (4)$$

By using the linearized thermal dynamics, we can describe the linear variation of the signal δV_S in terms of the linear response functions $\Theta_{\text{gs}}(T_s, T_g, \omega) \equiv \delta T_g / \delta T_s$ and $\Theta_{\text{sp}}(T_s, T_g, \omega) \equiv \delta T_s / \delta P_s$ (see Appendix A), obtaining

$$\mathcal{R} = \left(\frac{\partial V_S}{\partial T_s} + \frac{\partial V_S}{\partial T_g} \Theta_{\text{gs}}\right) \Theta_{\text{sp}}, \quad (5)$$

where we compute the derivatives of V_S by numerically computing V_S on a (T_s, T_g) grid [25]. We show the dependence of $\mathcal{R}$ on P_s in Fig. 2(b). $\mathcal{R}$ assumes very large negative values on the order of approximately -10^{11} V/W for low powers P_s up to approximately 1 fW. For powers P_s greater than approximately 1 fW, $\mathcal{R}$ inverts its sign and reaches relatively large positive values of approximately 10^9 V/W. This sign inversion is a nonlinear effect, described in Ref. [25], due to the fact that for large powers $V_S(P_s)$ is not a monotonic function [25]. Intuitively, for large P_s , the dominant contribution to changes in V_S is caused by the change in T_g and not by the change in T_s . This results in an inverted behavior, as an increase in P_s implies an increase in T_g , reducing the thermal gradient and thus the absolute value of V_S .

These zeroes in $\mathcal{R}$ have unfavorable implications for the noise figures of merit, as we will see below. However, the system can be used as a bolometer in both regimes as long as the power is far enough from the inversion point where $\mathcal{R} = 0$.

III. NOISE MODEL AND RESULTS

To assess device properties, modeling noise has the same importance as modeling the signal. While physical noise sources are the same at room temperature and at cryogenic temperature, dominant contributions dramatically change for temperatures lower than the Bloch-Grüneisen temperature [17,24,32,33]. As stated above, we focus on the open-circuit configuration.

Noise registered in the measurement comes from charge and energy fluctuations in the quasiparticle gases on both sides of the junction. Charge fluctuations are directly measured via the induced potential fluctuations δV_S , while the effect of energy fluctuations on steady-state temperatures can be described via Eq. (A1) if the variation is slow enough ($\omega \ll 1/\tau_{\text{tun}}$, where $\tau_{\text{tun}} = e/I \sim 10^{-9}$ s is the average time between electron tunneling events).

We can compare the effects of very different noise contributions by showing their NEP, an important bolometric figure of merit [13,17] defined as the input power δP_s required to generate a signal equal to the average value of the noise δV_S with a 1 Hz sampling interval. Following standard literature [17,24], we now describe how we estimate the NEP for each dominant noise contribution before computing the total NEP by quadrature sum of all contributions [34].

A. Johnson-Nyquist noise

Any resistive element with resistance R in a circuit induces an open-circuit voltage noise $\langle \delta V^2 \rangle = 4k_B RT \delta \omega$ [35]. With our geometry, the graphene layer induces a voltage noise δV_g equivalent to $\sqrt{\langle \delta V^2 \rangle} \sim 5.9 \times 10^{-10}$ V/ $\sqrt{\text{Hz}}$ at $T_g = 1$ K if we assume a graphene mobility μ_g of 10^3 cm²/V s and a carrier density n_0 of 10^{12} cm⁻².

However, this noise has to be summed incoherently with the noise from the amplifier circuit required to read V_S , which we can estimate to be $\delta V_{\text{amp}} \sim 10^{-9}$ V/ $\sqrt{\text{Hz}}$ [36]. Therefore, we only keep δV_{amp} for estimation purposes. We can then compute the corresponding NEP by using the definition of responsivity $\text{NEP}_{\text{amp}} = \delta V_{\text{amp}}/\mathcal{R}$ [17].

B. Phonon noise

The interaction between quasiparticles and phonons induces energy fluctuations in a quasiparticle gas [32,34,37]. These fluctuations happen in both the superconductor and the graphene layer, and they respectively depend on T_s and T_g . We assume a negligible Kapitza resistance; therefore, we consider the phonon gas to be at the thermal bath temperature T_b .

For the superconductor layer, we estimate the phonon noise power spectral density by using the equilibrium formula $\delta P_{\text{q-ph},s}^2 = 4k_B G_{\text{q-ph},s} T_s^2 \delta \omega$ [17,32,37], with the quasiparticle-phonon thermal conductivity defined as $G_{\text{q-ph},s} = \partial Q_{\text{s-ph}}/\partial T_s$. Even if the quasiparticles and the phonons are not at equilibrium when $T_s \neq T_b$, we believe this formula to be a cautious approximation because it corresponds to considering the phonon gas to be at the same temperature as the quasiparticle gas, therefore inducing a larger number of interactions and thus strongly overestimating noise. This noise induces power fluctuations $\delta P_{\text{q-ph},s}$ in the superconductor gas, and therefore it adds to the signal P_s , thus implying $\text{NEP}_{\text{q-ph},s} = \delta P_{\text{q-ph},s}$.

For the graphene layer, we estimate the phonon noise power spectral density in the clean limit for graphene $\delta P_{\text{q-ph},g}^2 = 8k_B A_g \Sigma_C (T_g^5 + T_b^5)$, where Σ_C is the electron-phonon coupling constant in the clean limit [24]. This noise acts as an oscillating δP_g in the second line of Eq. (A1). This means that it does not generate a δV_S oscillation with the same mechanism as a δP_s oscillation, and that we cannot equate the NEP to $\delta P_{\text{q-ph},g}$, and we need to compute the response function $\mathcal{R}_g \equiv \delta V_S/\delta P_g$ and translate the δV_S oscillation back to an equivalent input power fluctuation $\delta V_S/\mathcal{R} \equiv \text{NEP}$, which is the NEP by definition. After algebraic manipulation of Eq. (A1), we obtain

$$\mathcal{R}_g = \left[\frac{\partial V_S}{\partial T_g} \left(1 - \frac{\partial Q_{\text{GIS}}}{\partial T_g} \Theta_{\text{gs}} \right) - \frac{\partial V_S}{\partial T_s} \frac{\partial Q_{\text{GIS}}}{\partial T_g} \Theta_{\text{sp}} \right] \mathcal{K}_g, \quad (6)$$

where the expression for $\mathcal{K}_g(T_s, T_g, \omega)$ can be found as Eq. (A6) in Appendix A. We can then compute $\text{NEP}_{\text{q-ph},g} = (\mathcal{R}_g/\mathcal{R}) \delta P_g$.

C. Junction noise

The stochasticity of the tunneling process implies that the amount of charge and energy flowing through the junction fluctuates [32,38]. We estimate the contribution to the total NEP by using the fluctuations in current $\langle \delta I^2 \rangle$,

power $\langle \delta P_j^2 \rangle$, and their cross-correlation $\langle \delta P_j \delta I \rangle$ computed from the statistical fluctuations in the number of tunneling quasiparticles by weighing the fluctuations on charge and energies [17,24,32], with full expressions contained in Appendix B.

Similarly to $\text{NEP}_{\text{q-ph,g}}$, δP_j does not simply enter Eq. (A1) as a δP_s oscillation, but rather it acts as a simultaneous power oscillation on both sides $\delta P_g = -\delta P_s = \delta P_j$ because of energy conservation. We use algebra analogous to Eq. (6) to compute the response function $\mathcal{R}_{P_j} \equiv \delta V_S / \delta P_j$, obtaining

$$\mathcal{R}_{P_j} = -\frac{\partial V_S}{\partial T_s} \left(\frac{\partial Q_{\text{GIS}}}{\partial T_g} \mathcal{K}_g + 1 \right) \Theta_{\text{SP}} + \frac{\partial V_S}{\partial T_g} \left[\mathcal{K}_g - \Theta_{\text{gs}} \Theta_{\text{SP}} \left(\frac{\partial Q_{\text{GIS}}}{\partial T_g} \mathcal{K}_g + 1 \right) \right]. \quad (7)$$

While we compute the function for the case $P_g = 0, P_s \neq 0$, we can find analogous functions for different device configurations.

We estimate the open-circuit voltage noise from current fluctuations by using the standard formula $\delta V_I = (dI/dV)^{-1} \delta I \equiv R_d \delta I$ [17,32,38], where we use the derivative of Eq. (1). Considering that we can use the same rules for the cross-correlation [17], the final result for the NEP corresponding to junction noise is

$$\text{NEP}_j^2 = \frac{\mathcal{R}_{P_j}^2 \langle \delta P_j^2 \rangle + R_d^2 \langle \delta I^2 \rangle - 2\mathcal{R}_{P_j} R_d \langle \delta P_j \delta I \rangle}{\mathcal{R}^2}, \quad (8)$$

where the relative sign between the correlation and the squared noise contributions comes from the tunneling process properties [32].

In Fig. 3(a), we can see the results for every NEP contribution computed for our device setup with $T_b = 10$ mK. We can see that the junction contribution overwhelmingly dominates over every other noise source. This is mainly due to the large contribution of the $\langle \delta I^2 \rangle$ part, which is multiplied by a large dynamic resistance $R_d \sim 1$ G Ω . This large R_d comes from the junction still being away from the Ohmic part of the I - V curve at $V = V_S$ for our values of the Seebeck potential. The sharp peaks present in each contribution except $\text{NEP}_{\text{q-ph,s}}$ correspond with the $\mathcal{R} = 0$ point shown in Fig. 2(b).

In Fig. 3(b), we represent the dependence of the total NEP on input power P_s for different values of T_b . This is the most important figure of merit for a bolometer, and for our setup its most useful value is approximately 4×10^{-17} W/ $\sqrt{\text{Hz}}$ estimated for $P_s \sim 10^{-16}$ W and $T_b = 100$ mK, since it shows similar behavior to the total NEP at 10 mK. Another important feature is that, for intermediate powers 10^{-16} W $\lesssim P_s \lesssim 10^{-14}$ W, the higher bath temperature $T_b = 300$ mK shows a performance similar to the lower T_b , which is useful for further relaxing requirements on the cryogenic apparatus.

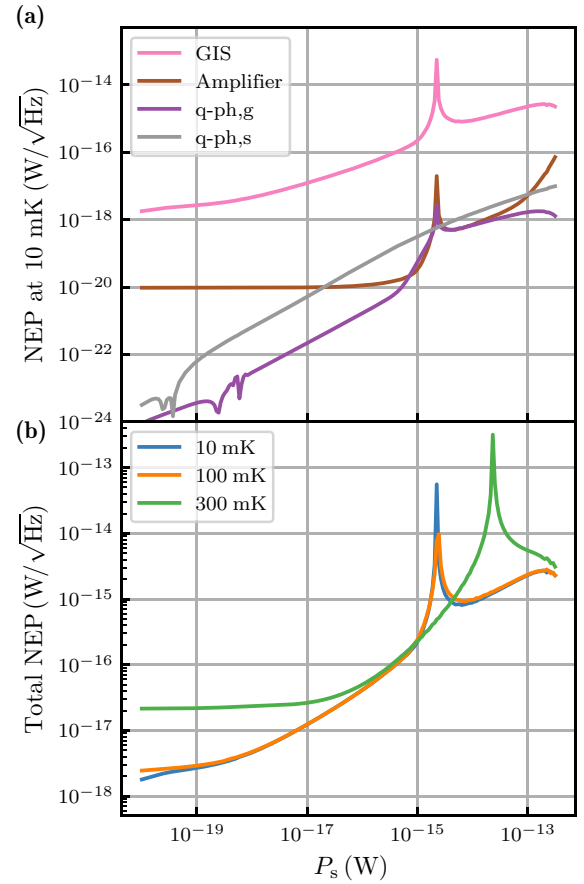


FIG. 3. Dependence of partial and total NEP on P_s . (a) NEP contribution at $T_b = 10$ mK for each noise source. Junction noise dominates by several orders of magnitude because of a large dynamical resistance R_d . The peak represents the zero of the responsivity $\mathcal{R}$. (b) Comparison of total NEP dependences on P_s for different T_b values. The best reasonable value for the NEP is 4×10^{-17} W/ $\sqrt{\text{Hz}}$ for $P_s \sim 10^{-16}$ W for $T_b = 100$ mK, but the choice of T_b could depend on the required input power for the device.

IV. CHARACTERISTIC TIMES

Characteristic times represent another important consideration for a bolometer design. The two most important timescales to represent the behavior of the bolometer are the response time and the integration time needed to obtain a useful measurement.

Since electrical response times are much faster, we estimate the bolometer response time via the thermal response time τ_{th} , the time needed by the system to adjust its temperature after a change in the power input. We imagine a response to a small variation δP over a large power input P_s , which makes the system reach a new steady state exponentially via Eq. (3), with a characteristic time

$$\tau_{\text{th}} = \frac{C_s(T_s)}{G_{\text{GIS}}(T_s, T_g) + G_{\text{q-ph,s}}(T_s)}, \quad (9)$$

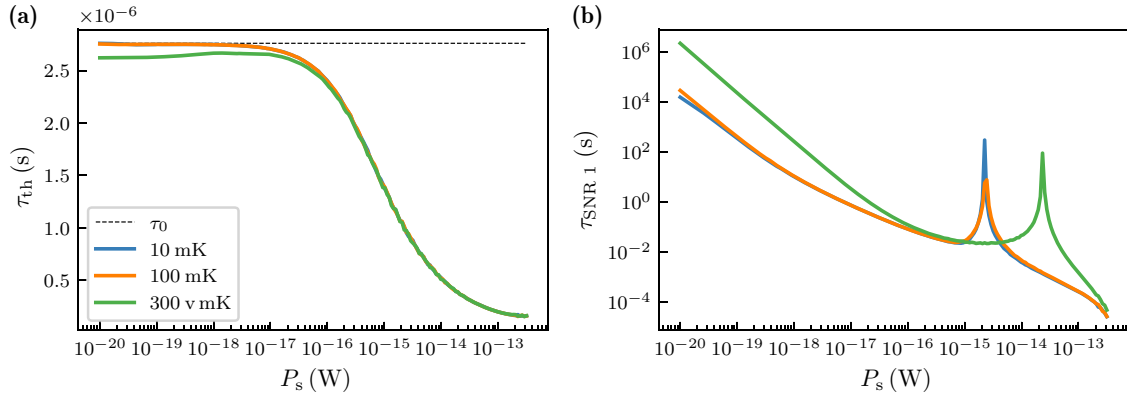


FIG. 4. Dependence of characteristic times on P_s . (a) Thermal response time τ_{th} for different T_b values. The curves for 10 and 100 mK are almost identical. For small P_s and T_b , $\tau_{th}(P_s) \sim \tau_0$ shows almost no temperature dependence (see the main text). τ_{th} decreases for larger P_s because of increased energy dissipation in the superconductor. (b) Integration time necessary to get a signal-to-noise-ratio equal to 1, $\tau_{SNR 1}$. While our bolometer shows very long integration times of approximately 10^3 s for low power, the integration time decreases with increasing P_s , decreasing to approximately 1 ms for $P_s \sim 10^{-13}$ W and $T_b = 100$ mK. Integration times are computed for signal variations corresponding to the entire signal.

where $G_i = \partial Q_i / \partial T_s$. This standard approximation [34, 37] is reasonable for estimation purposes, and it gives a correct physical picture for experiments where the signal variations are rather slow, such as the sky scans for CMB survey. We show the dependence of τ_{th} on P_s in Fig. 4(a). For $P_s < 10^{-17}$ W, τ_{th} is constant for $T_b = 10$ mK and 100 mK, then it decreases by an order of magnitude from $P_s \sim 10^{-17}$ W to $P_s \sim 10^{-13}$ W. For $T_b = 300$ mK, τ_{th} is almost constant. For $T_b = 300$ mK, τ_{th} shows the same behavior as for $T_b = 100$ mK, except for a different limit value at lower P_s , due to non-negligible quasiparticle dissipation in the superconductor (see Fig. 4 in Ref. [25]). We prove in Appendix C that for small P_s τ_{th} converges to a fixed value $\tau_0 = \mathcal{V}_s N_F e^2 R_T \sim 2.76 \mu s$, as long as $G_{q-ph,s} \ll G_{GIS} T_g \ll T_s$ and $T_s \ll \Delta_0 / k_B$. In Fig. 4 in Ref. [25], we can see that $Q_{s-ph,s}$ is very small for $T_s \lesssim 0.3$ K, and thus $G_{q-ph,s}$ is negligible except for $T_b = 300$ mK. We can see that τ_{th} reaches approximately 200 ns for $P_s \sim 10^{-13}$ W. The effect of $G_{q-ph,s}$ also inverts the dependence of τ_{th} on P_s , making it drop, while G_j alone would make it rise.

We define the integration time

$$\tau_{SNR 1} = \frac{1}{2} \left(\frac{NEP_{tot}(T_s(P_s), T_g(P_s))}{P_s} \right)^2 \quad (10)$$

as the time that a signal of fixed power P_s needs to be integrated to have a signal-to-noise ratio (SNR) of 1. Integration for a longer time reduces the frequency interval and thus reduces the absolute value of the noise power. In Fig. 4(b), we show the dependence of $\tau_{SNR 1}$ on P_s . Interestingly, $\tau_{SNR 1}$ becomes smaller with increasing P_s , down to approximately 100 μs for $P_s \sim 10^{-13}$ W. The $\mathcal{R} = 0$

peaks are also present in this plot, due to the NEP in the numerator.

V. CONCLUSIONS

In this article, we have proposed and numerically characterized a thermoelectric bolometer made of a graphene-insulator-superconductor junction. After assessing that we have better device performance by heating the superconductor side, we discovered that the device shows a large responsivity $|\mathcal{R}|$ of up to approximately 10^{11} V/W for $P_s \lesssim 1$ fW and $T_b = 100$ mK. Furthermore, our device design shows a NEP of approximately 4×10^{-17} W/ $\sqrt{Hz}$ for $P_s \sim 10^{-16}$ W and $T_b = 100$ mK, and a NEP $\sim 2 \times 10^{-15}$ W/ $\sqrt{Hz}$ for $P_s \sim 10^{-13}$ W. Even though our NEP is larger than the measured approximately 10^{-17} W/ $\sqrt{Hz}$ reported in the literature [9] and the requirement for future experiments, our detector is passive with performance almost comparable to that of active detectors. In the preparation of a large-array experiment, a smaller impact on the heat budget and cheaper fabrication could be more useful than a smaller NEP, which can, in principle, be compensated by longer integration times. Since this is a preliminary study, further optimization could also improve the noise performance of our device. Our detector shows response times τ_{th} down to approximately 200 ns for $P_s \sim 10^{-13}$ W, and integration times $\tau_{SNR 1}$ down to approximately 100 μs for $P_s \sim 10^{-13}$ W. These values are compatible with those from investigations of signals arising from experiments focused on interesting physical phenomena, such as the polarization of the CMB [7]. Thus, this technology might be implemented in experiments where the total power budget of the apparatus is limited, such as balloon or satellite-based cosmology experiments.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX A: RESPONSE FUNCTIONS AND LINEARIZATION OF THERMAL DYNAMICS

To obtain the response functions to small oscillations in input power, we start from the linearized version of the thermal dynamics equations

$$\begin{cases} i\omega C_s \delta T_s = - \left(\frac{\partial Q_{\text{GIS}}}{\partial T_s} + \frac{\partial Q_{\text{s-ph}}}{\partial T_s} \right) \delta T_s - \frac{\partial Q_{\text{GIS}}}{\partial T_g} \delta T_g + \delta P_s \\ i\omega C_g \delta T_g = \left(\frac{\partial Q_{\text{GIS}}}{\partial T_g} - \frac{\partial Q_{\text{g-ph}}}{\partial T_g} \right) \delta T_g + \frac{\partial Q_{\text{GIS}}}{\partial T_s} \delta T_s + \delta P_g. \end{cases} \quad (\text{A1})$$

To understand the response to a variation in δP_s , we algebraically solve the system first for δT_g , obtaining

$$\delta T_g = \Theta_{\text{gs}} \delta T_s, \quad (\text{A2})$$

and then we solve for δT_s , obtaining

$$\delta T_s = \Theta_{\text{sp}} \delta P_s, \quad (\text{A3})$$

with

$$\Theta_{\text{gs}}(T_s, T_g, \omega) \equiv \frac{\delta T_g}{\delta T_s} = \frac{\frac{\partial Q_{\text{GIS}}}{\partial T_s}}{i\omega C_g - \frac{\partial Q_{\text{GIS}}}{\partial T_g} + \frac{dQ_{\text{g-ph}}}{dT_g}}, \quad (\text{A4a})$$

$$\begin{aligned} \Theta_{\text{sp}}(T_s, T_g, \omega) &\equiv \frac{\delta T_s}{\delta P_s} \\ &= \frac{1}{i\omega C_s - \frac{\partial Q_{\text{GIS}}}{\partial T_g} \Theta_{\text{gs}}(T_s, T_g, \omega) - \frac{\partial Q_{\text{GIS}}}{\partial T_s} + \frac{dQ_{\text{s-ph}}}{dT_s}}. \end{aligned} \quad (\text{A4b})$$

The same set of equations could be solved differently to obtain the response to a small δP_g with a fixed P_s . In that case, we still solve first for δT_g , but this time we obtain

$$\delta T_g = \Theta_{\text{gs}} \delta T_s + \mathcal{K}_g \delta P_g, \quad (\text{A5})$$

with

$$\mathcal{K}_g(T_s, T_g, \omega) = \frac{1}{i\omega C_g - \frac{\partial Q_{\text{GIS}}}{\partial T_g} + \frac{dQ_{\text{g-ph}}}{dT_g}}. \quad (\text{A6})$$

The second part is the same as before, and it allows us to find the response functions to power fluctuations in the graphene quasiparticle gas $\mathcal{R}_g$ and to the tunneling-induced power fluctuations in the junction $\mathcal{R}_{P_j}$, as described in the main text.

APPENDIX B: NOISE EQUIVALENT POWER EXPRESSIONS FOR THE JUNCTION

In the main text, we use the expectation values of the fluctuations in current $\langle \delta I^2 \rangle$, power transmitted across the junction $\langle \delta P_j^2 \rangle$, and their cross-correlation $\langle \delta P_j \delta I \rangle$. Following Ref. [24], we compute these quantities by computing the expectation values of the second moment of the relevant transported quantity (charge, energy, or both) on the energy-dependent tunneling probabilities. The full expressions are

$$\begin{aligned} \langle \delta P_j^2 \rangle &= \frac{2}{e^2 R_T} \int dE (E - eV)^2 \rho_g(E - E_F - eV) \rho_s(E, T_s) \\ &\quad \times [f(E, T_s) + f(E - eV, T_g) \\ &\quad - 2f(E, T_s)f(E - eV, T_g)], \end{aligned} \quad (\text{B1})$$

$$\begin{aligned} \langle \delta I^2 \rangle &= \frac{2}{R_T} \int dE \rho_g(E - E_F - eV) \rho_s(E, T_s) \\ &\quad \times [f(E, T_s) + f(E - eV, T_g) \\ &\quad - 2f(E, T_s)f(E - eV, T_g)], \end{aligned} \quad (\text{B2})$$

$$\begin{aligned} \langle \delta P_j \delta I \rangle &= \frac{2}{e R_T} \int dE (E - eV) \rho_g(E - E_F - eV) \rho_s(E, T_s) \\ &\quad \times [f(E, T_s) + f(E - eV, T_g) \\ &\quad - 2f(E, T_s)f(E - eV, T_g)]. \end{aligned} \quad (\text{B3})$$

APPENDIX C: DEMONSTRATION OF CONSTANT THERMAL RESPONSE TIME FOR SMALL INPUT POWERS

In Fig. 4(b), we can see that for low input powers P_s and low bath temperatures T_b the response time is almost constant. The thermal response time of a bolometer is usually

defined as [34]

$$\tau_{\text{th}} = \frac{C_{\text{th}}}{G_{\text{th}}}, \quad (\text{C1})$$

where C_{th} is the thermal capacitance of the active region and G_{th} is the thermal conductance between the charge carriers in the active region and the thermal bath. In the linear regime, C_{th} and G_{th} are computed at the bath temperature; therefore, they are constant for every input power. However, as shown in Ref. [25], the linear regime breaks for power P_s much smaller than approximately 10^{-17} W, where we still observe an almost constant τ_{th} . Here we prove that the constant value of τ_{th} holds even in the nonlinear regime, as long as T_s and T_g are small enough.

When P_s is small, the bottleneck for heat conduction is the junction G_{GIS} , as phonon-induced dissipation in the superconductor is negligible and phonon-induced dissipation in graphene is strong enough to keep $T_g \sim T_b$. Therefore, the thermal capacitance of the superconductor acts as the thermal capacitance of the entire system, and the Seebeck potential is negligible $V_s \sim 0$. From the definitions we obtain

$$C_s(T_s) = T_s \frac{dS_s}{dT_s} = -2\mathcal{V}_s N_F k_B T_s \frac{d}{dT_s} \times \int_{-\infty}^{+\infty} dE \rho_s(E) f(E, T_s) \ln f(E, T_s), \quad (\text{C2})$$

$$G_{\text{GIS}}(T_s, T_g) = \frac{\partial Q_{\text{GIS}}}{\partial T_s} = \frac{1}{e^2 R_T} \frac{\partial}{\partial T_s} \times \int dE E \rho_g(E - E_F) \rho_s(E) [f(E, T_g) - f(E, T_s)], \quad (\text{C3})$$

where R_T is the tunneling resistance, $f(E, T)$ is the Fermi function, ρ_s is the DOS of the superconductor and ρ_g is the DOS of graphene defined by Eqs. (2c), (4), and (8) in Ref. [25], $N_F = 2.15 \times 10^{47} \text{ J}^{-1} \text{ m}^{-3}$ is the superconductor DOS at the Fermi level in the metallic state, and $\mathcal{V}_s$ is the superconductor volume. Since we consider doped graphene with electron density $n_0 = 1 \times 10^{12} \text{ cm}^{-2}$, we can assume the relative graphene DOS $\rho_g(E - E_F) = 1$ to be approximately constant over the relevant energy interval, and we consider T_s to be low enough to not impact the superconductor DOS ρ_s .

If we focus on the T_s -dependent part of $C_s(T_s)$, we can rewrite Eq. (C2) as

$$T_s \frac{d}{dT_s} \int dE \rho_s(E) f(E, T_s) \ln f(E, T_s) = T_s \int dE \frac{\partial \rho_s}{\partial T_s} f \ln f + \rho_s \frac{\partial f}{\partial T_s} \ln f + \rho_s f \frac{\partial \ln f}{\partial T_s}. \quad (\text{C4})$$

The last term in Eq. (C4) is zero because $\rho_s(E)$ is even in E and $\partial f / \partial T_s$ is odd in E . In the $T_s \rightarrow 0$ limit ($T_s \ll \Delta_0 / k_B$),

we can recast the second term as

$$\lim_{T_s \rightarrow 0} \ln f = \begin{cases} -\frac{E}{k_B T_s}, & E > 0, \\ 0, & E < 0, \end{cases} \quad (\text{C5})$$

obtaining

$$\lim_{T_s \rightarrow 0} C_s(T_s) = 2\mathcal{V}_s N_F \int_0^{+\infty} dE E \left(\frac{\partial \rho_s}{\partial T_s} f + \rho_s \frac{\partial f}{\partial T_s} \right). \quad (\text{C6})$$

We now focus on the T_s -dependent part of G_{GIS} in Eq. (C3),

$$\begin{aligned} & \frac{\partial}{\partial T_s} \int_{-\infty}^{+\infty} dE E \rho_s(E) [f(E, T_s) - f(E, T_g)] \\ &= \int_{-\infty}^{+\infty} dE E \left[\frac{\partial \rho_s}{\partial T_s} (f(E, T_s) - f(E, T_g)) + \rho_s \frac{\partial f}{\partial T_s} \right]. \end{aligned} \quad (\text{C7})$$

The second term is even in E , and its integral between $-\infty$ and $+\infty$ is just 2 times its integral between 0 and $+\infty$. If $T_g \ll T_s$,

$$\lim_{T_g \ll T_s} f(E, T_s) - f(E, T_g) = \begin{cases} f(E, T_s), & E > 0, \\ f(E, T_s) - 1, & E < 0. \end{cases} \quad (\text{C9})$$

We can recast the negative energy part of the integral by using the variable transformation $E \rightarrow -E$ and the fact that $\partial \rho_s / \partial T_s$ is an even function of E :

$$\int_{-\infty}^0 dE E \frac{\partial \rho_s}{\partial T_s} (f(E, T_s) - 1) = \int_0^{+\infty} dE E \frac{\partial \rho_s}{\partial T_s} f(E, T_s), \quad (\text{C10})$$

which is equal to the positive energy part. Therefore,

$$\lim_{T_g \ll T_s} G_j = \frac{2}{e^2 R_T} \int_0^{+\infty} dE E \left(\frac{\partial \rho_s}{\partial T_s} f + \rho_s \frac{\partial f}{\partial T_s} \right). \quad (\text{C11})$$

The T_s -dependent parts of C_s and G_j are equal, so if $T_s \rightarrow 0$ and $T_g \ll T_s$, τ_{th} remains constant at the limiting value

$$\lim_{T_g \ll T_s \rightarrow 0} \frac{C_s}{G_j} = \mathcal{V}_s N_F e^2 R_T \sim 2.76 \text{ } \mu\text{s}. \quad (\text{C12})$$

In our setup, T_s and T_g follow the conditions even for $P_s > 10^{-17}$ W, but the thermal conductivity induced by quasiparticle-phonon dissipation in the superconductor for $P_s > 10^{-17}$ W is not negligible anymore, inducing the change in τ_{th} described in the main text.

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