Far-infrared electrodynamics of thin superconducting NbN film in magnetic fields

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Abstract
We studied a thin superconducting NbN film in magnetic fields up to 8 T above the zero-temperature limit by means of time-domain terahertz and scanning tunneling spectroscopies in order to understand the vortex response. Scanning tunneling spectroscopy was used to determine the optical gap and the upper critical field of the sample. The values obtained were subsequently used to fit the terahertz complex conductivity spectra in the magnetic field in the Faraday geometry above the zero-temperature limit. These spectra are best described in terms of the Coffey–Clem self-consistent solution of a modified London equation in the flux creep regime.

Keywords: superconductivity, thin films, scanning tunneling spectroscopy, time-domain terahertz spectroscopy

(Some figures may appear in colour only in the online journal)

1. Introduction
Terahertz-range electrodynamics of superconductors is governed by the response of Cooper pairs and thermally activated quasiparticles [1–4]. The key parameter is the optical gap $2\Delta$, which represents the energy necessary for breaking Cooper pairs [5]. For classical BCS-like superconductors, the value of the optical gap ranges from tenths to units of meV, which corresponds to the high-frequency part of the microwave range and to the terahertz (THz) region.

In a magnetic field $B$ lying between the lower and upper critical fields ($B_{c1} < B < B_{c2}$), type-II superconductors, such as NbN, enter into the so-called Abrikosov or mixed state. In this state, the magnetic field penetrates the superconductor through cylindrical regions carrying a quantized magnetic flux $\Phi_0$, so-called vortices: the normal state is locally restored in vortex cores. The vortex lattice has a large impact on the high-frequency electrodynamical response of superconductors not only because the material becomes inhomogeneous, but also because a vortex oscillatory motion inevitably contributes to both the inductive and the dissipative parts of the electrodynamic response [6–8]. While the vortex state has been studied quite thoroughly in the microwave range [9–12], there are only few studies in the far-infrared range lying in the vicinity of the optical gap [13, 14]. The conclusions of these studies, however, do not provide a simple unambiguous picture: Ikebe et al [13] explained their experiments on classical NbN superconductor by the Coffey–Clem model within the dissipative flux-flow regime, while Xi et al [14] argued that the vortices are strongly pinned and that the contribution from the vortex motion can be neglected.
The aim of this paper is to clarify the issue by studying a superconducting NbN film in the Faraday geometry in a broad temperature range between \( T = 0 \) K and a temperature slightly above the critical temperature \( T_c \). A number of important parameters of the NbN film studied were determined by supplemental scanning tunneling spectroscopy (STS) experiments, which enabled us to test various theoretical models with a minimum of free parameters.

2. Experiment

NbN was chosen as a typical representative of classical BCS type-II superconductors. Our sample was a 11.5 nm thick NbN film deposited on a highly resistive Si substrate. The nominal critical temperature determined by DC conductivity measurements was \( T_c = 11.5 \) K.

A custom-made spectrometer was used for time-domain terahertz spectroscopy (TDS) measurements in the transmission geometry arrangement. Broadband THz pulses were generated using a Ti:sapphire femtosecond laser and a commercial large-area semiconductor interdigitated emitter (TeraSED, GigaOptics). The sample was placed in an Oxford Instruments Spectromag He-bath cryostat with mylar windows (TeraSED, GigaOptics). The sample was placed in an Oxford Instrument cryostat with a cold He bath. The sample was several degrees higher than the readout on the sample. The sample was then scanned using the NbN sample. A bias voltage was applied to the tip, while the sample was grounded. The Au tip enabled the formation of an N–I–S tunnel junction, where N represents the gold tip (normal state), I is the insulating barrier of the vacuum, respectively the surface oxides, and S stands for the superconducting NbN sample.

The tunneling spectrum, i.e. the differential conductance versus the voltage dependence observed by STS, represents a convolution of the local density of states of both electrodes comprising the tunneling junction. Since the normal tip features a constant density of states near the Fermi energy, the differential conductance \( dI/dV \) versus voltage of an N–I–S junction reflects the local energy-dependent superconducting density of states \( N_S(E) \) smeared by \( k_B T \) in energy at temperature \( T \), where \( k_B \) is the Boltzmann constant:

\[
\frac{dI}{dV}(V) \propto \int_{-\infty}^{\infty} N_S(E) \frac{df}{dE}(E + eV) dE, \tag{2}
\]

where \( f \) is the Fermi–Dirac distribution. The tunneling probability which also enters into the expression of the tunneling current is taken as a constant at low voltages where the superconducting energy gap is scanned. For a BCS superconductor, \( N_S(E) = E/\sqrt{E^2 - \Delta^2} \), where \( \Delta \) is the superconducting energy gap. Consequently, in the low-temperature limit \( (k_B T \ll \Delta) \), the differential conductance measures the superconducting density of states \( N_S \) directly \[19\].

3. Results

First, our experimental results in zero magnetic field are evaluated, then our efforts are extended to include magnetic field measurements in the Faraday geometry.

3.1. Zero magnetic field

STM measurements have revealed that an oxidized layer covers the whole surface of our NbN film, thus preventing us from a detailed topographic investigation of the sample surface. Semenov et al \[20\] estimated the thickness of oxidized...
Figure 1. Differential tunneling conductance of the NbN–Au tunnel junction measured at several temperatures between 0.7 and 12.5 K. The inset shows the temperature dependence of the energy gap \( \Delta(T) \) in a zero magnetic field together with the prediction of the BCS model (full line).

The temperature dependence of the gap was determined by fitting the spectra by equation (2) taking into account the tunneling conductance taken at the lowest temperature. The value of the superconducting gap \( 2\Delta(0) \) corresponds roughly to the distance of the peaks in the tunneling conductance taken at the lowest temperature. The temperature dependence of the gap was determined by fitting the spectra by equation (2) taking into account the thermal broadening. Moreover, in the BCS density of states expression [21], this enables us to account for the inhomogeneity of the sample. From the fit at 0.7 K we obtained a relatively small value of \( \Gamma = 0.12\Delta(0) \), which was then temperature independent. The temperature dependence of the superconducting gap obtained is shown in the inset of figure 1, together with the prediction of the BCS theory [19] adjusted for \( 2\Delta(0)/k_B T_c = 4.25 \), and a reasonable agreement is found.

The superconducting energy gap \( \Delta(0) \) and the local transition temperature \( T_c \) were checked at several locations on the sample surface. The local values of the gap \( \Delta(0) \) estimated from the fits of the low-temperature spectra to equation (2) varied between 2.1 and 2.3 meV. The local \( T_c \) was independently determined from the temperature dependence of the energy gap \( \Delta(T) \) and from its extrapolation to zero. The resulting \( T_c \) varied between 11.5 and 12 K. These results point to a moderate inhomogeneity of the NbN film.

Since the thickness of our sample corresponds only to a few coherence lengths, we conclude that this inhomogeneity is rather in plane than along the depths of the film. Indeed, the coherence length \( \xi(0) = \sqrt{\Phi_0/2\pi B_{c2}(0)} \) estimated from the upper critical field \( B_{c2}(0) \approx 26 \text{ T} \) (see below) amounts to 3.5 nm. Thus, our STS measurements determine a ratio of \( 2\Delta(0)/k_B T_c = 4.25-4.5 \) in our NbN films, suggesting a strong coupling superconductivity.

For steady-state TDTS experiments, a bare 404 \( \mu \text{m} \) thick highly resistive Si substrate was used as a reference. First, the normal-state properties of the NbN film were determined just above the superconducting phase transition. Using the Drude model \( \sigma_n(\omega) = \sigma_n(0) / (1 - i\omega \tau) \), the value of DC conductivity \( \sigma_n(0) = 0.45 \mu\Omega^{-1} \text{m}^{-1} \) was obtained, while the relaxation time \( \tau \) could not be reliably determined. We estimated \( \tau < 15 \text{ fs} \); the best fit yielded \( \tau = 8 \text{ fs} \) (1/(2\(\pi\)\(\tau\) = 20 THz).

As pointed out above, the temperature sensor in the cryostat was not in a direct contact with the sample, and some temperature gradient appeared. This did not influence the measurements in the normal state, since the normal-state properties are only weakly temperature dependent. Below \( T_c \), we fitted the measured complex conductivity at each temperature by the Zimmerman model [22], with the temperature as the only free parameter (see figure 2). We observe in figure 2 that \( \sigma_2 \) exhibits a typical \( 1/\omega \) dependence, and that it decreases with increasing temperature, since the density of the condensate decreases. The sharp-edge feature in \( \sigma_1 \) is a characteristic fingerprint of the superconducting gap; it becomes less pronounced at higher temperatures due to an increasing number of quasiparticles. Close to \( T_c \),
\[ \sigma_1 \] continuously transforms into \[ \sigma_{1n} \]. Note that a careful evaluation of the experimental errors (error bars in figure 2) shows that \[ \sigma_2 \] is determined more precisely than \[ \sigma_1 \], and that especially below the superconducting gap (where \[ \sigma_1 \] acquires quite small values) the real part of the conductivity suffers from large relative uncertainties. This is an important statement, which will be used for fitting the experimental spectra obtained in the magnetic field.

### 3.2. Non-zero magnetic field

Tunneling measurements in magnetic fields up to 8 T were performed at different temperatures. A typical example of the measured data is shown in figure 3 with the spectra taken at 6 K in fields from 0 to 8 T with a 0.5 T step. As can be seen from the position of the peaks, the superconducting gap is only slightly altered by the magnetic field, and the normalized conductance at zero bias increases with increasing magnetic field. The Au STM tip was pressed against the sample surface to overcome the surface oxide layers. We deduced from imprints observed in an optical microscope after the STM experiments that the STM junction was of submicron size, in contrast to the nanometer sizes which are typical for usual STM experiments. In such a situation when a magnetic field \( B > B_{c1} \) is applied perpendicularly to the junction’s interface on the NbN film, not only the film but also the junction area is in the mixed state comprising the Abrikosov vortices. Then, the measured tunneling conductance corresponds to an in-plane average (over the junction’s area) of the local densities of states. The zero-bias tunneling conductance is a measure of the averaged quasiparticle density of states at the Fermi level. To a first approximation, the vortex cores represent the normal-state areas while the rest of the junction is superconducting, with a fully developed superconducting order parameter.

The increasing zero-bias tunneling conductance is thus proportional to the increasing volume fraction of vortex cores, which shows a linear dependence on the applied magnetic field \( (f_1 = V_{ch}/V \propto B/B_{c2}(T)) \). In the paper of Samuely et al [23] it was shown that the measurement of normalized zero-bias tunneling conductance as a function of the field strength is a very sensitive method to determine the upper critical magnetic field \( B_{c2}(T) \) for which the sample enters the normal state.

A substantial advantage of this method is that, even if the value of \( B_{c2}(T) \) exceeds the maximum attainable field, the upper critical field at a given temperature can be reliably determined by a linear extrapolation of the zero-bias tunneling conductance to its normal-state value.

The linear extrapolation of the normalized zero-bias tunneling conductance versus \( B \) is shown in the inset of figure 4. In this way, \( B_{c2}(T) \) was determined for several temperatures, and the results of this analysis are shown in the main part of the figure. Although they were obtained only in a limited temperature range, the \( B_{c2} \) values reveal a linear increase upon decreasing temperature below \( T_c \), with a tendency to saturation below 0.5 \( T/T_c \). Due to a lower number of data points, two different theoretical model \( B_{c2}(T) \) curves were drawn in order to extrapolate the zero-temperature value of the upper critical field. The classical Werthammer–Helfand–Hohenberg (WHH) dependence [24] is shown by the solid line yielding a correct fit. The extrapolated value \( B_{c2}(0) = 26 \) T is found. Another frequently used model developed by Tinkham [19], i.e. \( B_{c2} = B_{c2}(0)(1-t^2)/(1 + t^2) \), where \( t = T/T_c(B = 0) \), provides a fit of almost the same quality (dashed line) with a slightly higher value: \( B_{c2}(0) = 28 \) T. With these \( B_{c2}(0) \) estimates one derives the zero-temperature coherence length \( \xi(0) = 3.5 \pm 0.1 \) nm.

We measured THz complex conductivity spectra in the Faraday geometry (\( B \) perpendicular to the film) for magnetic fields up to 7 T and at several temperatures starting from 7.8 K up to \( T_c \). Typical results are shown in figure 5: at other temperatures, the spectra are qualitatively the same. The values of \( \sigma_2(\omega) \) in the THz range slowly decrease with magnetic field, while the frequency dependence of \( \sigma_1 \) is more complex. At low THz frequencies, \( \sigma_1(\omega) \) grows with increasing magnetic field and exceeds the normal-state value \( \sigma_{1n}(0) \). As the upper critical field is approached, \( \sigma_1 \) starts to decrease, and it reaches the normal-state value \( \sigma_{1n}(0) \) at \( B_{c2}(T) \). This is difficult to
temperatures with a substantial number of quasiparticles.

The generalization of this approach for finite frequencies below the superconducting gap. However, the applicability of the Coffey–Clem model is limited to temperatures depending on the temperature and the external magnetic field; a mixture of a normal fluid and a superfluid, whose fractions are caused by vortex dynamics are described by the Coffey–Clem lattice (a random distribution is assumed instead). Effects linked to the regular arrangement of vortices in a hexagonal lattice (a self-consistent way, and one finds that

\[
\tilde{\sigma}_{n} = \frac{2f_n\tilde{\sigma}_s (\tilde{\sigma}_n - \tilde{\sigma}_s)}{(1 - f_n)(\tilde{\sigma}_n - \tilde{\sigma}_s) + 2\tilde{\sigma}_s + \tilde{\sigma}_s},
\]

As mentioned above, the superconducting film is considered as a system consisting of a superconducting matrix and cylindrical inclusions of the normal-state material representing vortex cores [25]. The radius of the vortex core is defined by the coherence length \(\xi\), i.e., it is of the order of a few nanometers; the inter-vortex distance is usually of the same order for fields applied in our experiment. The typical wavelengths of the far-infrared radiation (hundreds of micrometers) are much larger; therefore, the electromagnetic radiation cannot sense individual vortices, and an effective complex conductivity can be used to describe the macroscopic properties of the system. The upper frequency limit of the effective medium models is given merely by the size of the vortices.

The Maxwell–Garnet theory [26] (MGT) is suitable for dilute systems of inclusions (vortices) with a percolated matrix (superconducting state), and in the case of cylindrical inclusions it takes the form [27]:

\[
\tilde{\varepsilon}_{MG} = \frac{1 + f_n\tilde{\varepsilon}_s (\tilde{\varepsilon}_n - \tilde{\varepsilon}_s)}{1 - f_n(\tilde{\varepsilon}_n - \tilde{\varepsilon}_s) + 2\tilde{\varepsilon}_s + \tilde{\varepsilon}_s},
\]

\[
\tilde{\sigma}_{MG} = \frac{2f_n\tilde{\sigma}_s (\tilde{\sigma}_n - \tilde{\sigma}_s)}{(1 - f_n)(\tilde{\sigma}_n - \tilde{\sigma}_s) + 2\tilde{\sigma}_s + \tilde{\sigma}_s},
\]

where \(f_n = V_n/V\) is the volume fraction of vortex cores, and \(\tilde{\sigma}_s\) and \(\tilde{\sigma}_n\) are the complex conductivities of the superconducting and normal state, respectively. Recently, Rychetsky [28] argued that the MGT formula holds even for high concentrations of inclusions as long as the matrix is percolated.

The Bruggeman approach [29] assumes that both the superconducting and the vortex-core components are surrounded by an effective medium with effective permittivity \(\tilde{\varepsilon}_B\). Local field and effective conductivity values are obtained in a self-consistent way, and one finds that [27]

\[
\tilde{\sigma}_B = \frac{1}{2} \left( \beta \pm \sqrt{\beta^2 + 4\tilde{\varepsilon}_n\tilde{\varepsilon}_s} \right).
\]

Here, \(\beta = (1 - 2f_n)(\tilde{\varepsilon}_n - \tilde{\varepsilon}_s)\), and the sign is chosen so that a physically relevant solution \(\varepsilon_{1,2} \geq 0\) is obtained.

4.2. Coffey–Clem model

Coffey and Clem [6] solved the modified London equation self-consistently in the presence of vortices and found that the vortex motion can modify both the dissipative and the inductive parts of the complex conductivity. Explicitly, their model can be expressed as

\[
\tilde{\sigma}_{CC} = \frac{(1 - f)\tilde{\sigma}_s + f\tilde{\sigma}_n}{1 - (1 - f)\tilde{\sigma}_s / \tilde{\sigma}_{vd}}.
\]

where \(f = 1 - (1 - T^4/T_c^4)(1 - B/B_{c2}(T))\) is the normal-fluid fraction and \(\tilde{\sigma}_{vd}\) is the conductivity induced by the vortex dynamics [9]:

\[
\tilde{\sigma}_{vd} = \frac{n_s}{b} \left( \frac{1 + i\omega/\omega_0}{\epsilon + i\omega/\omega_0} \right),
\]

where the dimensionless parameter \(\epsilon\) varies between 0 (at zero temperature) and 1 (at \(T_c\)) and describes effects of
the flux creep, \( \omega_0 \) is an effective depinning frequency, and \( b = B/B_{c2}(T) \). Depending on the temperature, on the magnetic field, and on the ratio \( \omega_0/\omega_c \), three different regimes emerge [30]: a flux-creep regime (where the flux creep dominates and the response is mainly dissipative), a pinning regime (where the inductive response prevails), and a flux-flow regime (where vortices move collectively and the dissipative response dominates).

5. Discussion

5.1. Sources of errors

An important source of errors in thin-film measurements comes from the uncertainty in the substrate thickness, which significantly influences the phase of the transmitted THz wave [17]. The error bars shown in figure 2 take into account two effects: the statistical error in the time-domain THz wave form calculated from individual data accumulations and a systematic error due to a substrate thickness uncertainty of 1 \( \mu \text{m} \); they can be considered as typical for all other data shown in this paper.

The oxidized layer covering the NbN film influences the evaluation of the complex conductivity. On the one hand, the presence of an ultrathin niobium oxide layer does not have any influence on the THz transmission of the sample. On the other hand, if such a layer develops, the superconducting film thickness is necessarily reduced by \( \Delta d_{\text{lim}} \). This is a source of error in the evaluated conductivity, given as \( \Delta \sigma_{sc} = (\Delta d_{\text{lim}}/d_{\text{lim}}) \sigma_{sc} \). In our case, it may be responsible for an error of up to 9%. However, this error only rescales \( \sigma_{n}(0) \), so the qualitative analysis is not affected. Nevertheless, in principle this error should be added to the error bars shown in figure 2.

The in-plane inhomogeneity leading to the observed differences in the values of \( T_c \) can cause some minor variations of the complex THz conductivity. So long as the sample temperature is sufficiently far from \( T_c \), the inhomogeneity will however not play a significant role.

Finally, small systematic oscillations observed in the spectra of figure 2 are probably due to a very small error in the instrumental function (parasitic reflection in the waveform), which can lead to a visible Fabry–Perot-like effect in the instrumental function (parasitic reflection in the waveform), which can lead to a visible Fabry–Perot-like effect in the conductivity of a thin film.

In zero magnetic field, the value of the superconducting gap was determined from the normalized tunneling conductance. The ratio \( 2\Delta(0)/k_B T_c \) was observed to vary from 4.25 to 4.5 along the surface of the sample; these values are slightly higher than what is typically observed [4, 13, 31–33]; nevertheless, even higher values have also been found [34]. These values suggest that NbN is a strong-coupling superconductor and that, in general, the high-frequency (THz) response should be treated within the framework of Eliashberg formulas following Nam [35, 36]. The dominant effect of strong coupling is a uniform decrease of \( \sigma_2(\omega) \), while the effect on \( \sigma_1(\omega) \) is only minor. Nevertheless, based on previous good experience with analysis of similar data without accounting for the strong-coupling effects [4, 31, 34, 37, 38], we believe that these effects are not substantial, and that the BCS-theory-based Zimmermann model [22] describes the THz properties of NbN films adequately.

5.2. Modeling the experimental data

The complex conductivity of the normal state of our NbN sample is well described by the Drude model: it is almost purely real and frequency independent in the observed range; the relaxation time \( \tau \) thus reaches a value of a few femtoseconds only, and its precise value cannot be established reliably. Below \( T_c \), and in the absence of a magnetic field, the experimental data are well described by the Zimmermann model [22]: see figure 2. The two-fluid model [19] fails to describe the features connected with the gap which are clearly observed in the real part of the conductivity spectra; by contrast, it describes the imaginary part fairly well; see the zero-field data shown in figures 6 and 7.

In the analysis of the transmission measurements in a magnetic field, we compare our experimental results with those from the theoretical models presented in section 4. Almost all the parameters of these models were determined using STS or from the zero-field transmission measurements. In the frame of the Coffey–Clem model, only the creep parameter \( \epsilon \) and the effective depinning frequency \( \omega_0 \) are free parameters.

Let us start with an application of the effective medium theories, since the local field effects (almost negligible in a microwave range where \( |\sigma_1| \gg |\sigma_2| \)) become important at terahertz frequencies. In figure 5, the predictions by the Maxwell–Garnett and the Bruggeman theories are compared with experimental data for two different temperatures. For \( T = 9.4 \text{ K} \) and \( B \) varying from 0 to 7 T, the volume fraction of vortex cores \( f_n = B/B_{c2}(T) \) spans almost over the entire interval from 0 to 1 (normal state), providing thus an excellent
opportunity to test these theories. On the one hand, both theories give similar results and agree with experimental data for the frequencies above the gap and also for low magnetic fields. On the other hand, at low frequencies and high magnetic fields, both theories fail to describe the experimental values of $\sigma_2(\omega)$. Furthermore, in the frame of the Maxwell–Garnett theory, the real part of $\sigma_1$ is limited by $\sigma_\ell(0)$, which is in disagreement with the experimental data. In our analysis, we neglected the fact that the conductivity of the superconducting matrix $\tilde{\sigma}$ is modified by the magnetic field $[39, 40]$. Xi et al. [14] took this effect into account, but their numerical calculations lead to a further decrease in $\sigma_2$, which disagrees with our experimental data. We thus concluded that this effect cannot explain the discrepancies between the predictions of the effective medium models and the experimental data.

We agree with the argument of Xi et al. [14] that the topology assumed by the Bruggeman theory does not match that of the vortex state. From this point of view one may expect that the MGT should be more suitable. In our previous study [41], however, we found that our experimental data were both quantitatively and qualitatively better described by the Bruggeman theory. This theory accidentally mimics the effect of vortex dynamics at low frequencies and high magnetic fields, as we can assume on the basis of the present study.

In figures 6 and 7, the low-frequency data ($\hbar \omega < 2\Delta$) are compared with results from the Coffey–Clem model, equation (5). Since this model is not adequate for higher frequencies, the conductivity at higher frequencies is not shown. Due to the weak dependence of the complex conductivity on free parameters, $\epsilon$ and $\omega_0$, it cannot be reliably determined from our data. Therefore we compare our data with two limiting cases—the flux-flow regime (where $\omega \gg \omega_0$ and the flux creep can be neglected; see panel A) and the flux-creep regime (where $\omega \ll \omega_0$ and the flux creep dominates; see panel B). Our data (both the real and the imaginary parts of $\tilde{\sigma}_c$) suggest that flux flow does not occur, which is in disagreement with previous observations [10, 13]. However, the Coffey–Clem model in the case of flux creep provides an excellent description of $\sigma_2$ and a reasonable description of $\sigma_1$.

6. Summary

We have reported properties of a thin NbN film deposited on a high-resistivity Si substrate, obtained by time-domain THz and scanning tunneling spectroscopies. STS revealed the presence of a continuous oxide layer on the top of the film. We found that the optical gap $2\Delta(0)$ varies between 2.1 and 2.3 meV, and that the critical temperature $T_c$ varies between 11.5 and 12 K. This indicates a strong coupling superconductivity, with $2\Delta(0)/k_B T_c = 4.25–4.5$.

The values of the upper critical magnetic field $B_{c2}$ were determined by a linear extrapolation of the normalized zero-bias tunneling conductance as a function of the field strength. Its temperature dependence is well described by the WHH model, with $B_{c2}(0) = 26$ T. The measured zero-field THz complex conductivity is well described by the Zimmermann model [22].

For low magnetic fields, we found that both the Maxwell–Garnett and Bruggeman theories give similar results and agree with the experimental data. At high magnetic fields, however, they both fail to describe the dissipation at low frequencies and the inductive response, $\sigma_2$. The obvious reason is that they do not account for the effects of vortex dynamics, which are particularly important at low THz frequencies. Fortunately, the Coffey–Clem two-fluid approach, which neglects the pair-breaking process for frequencies above the optical gap, can be used for low THz frequencies. Our experimental results were compared with those from the Coffey–Clem model assuming either the flux-flow or the flux-creep regime. Based on the excellent description of $\sigma_2$ and the reasonable description of $\sigma_1$, we concluded that the Coffey–Clem model in the flux-creep regime best describes the THz properties of our NbN sample.

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