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LETTER TO THE EDITOR

Phonon escape from electrically heated metal films on silicon

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Abstract. Electric field heating of electrons in thin metal films deposited on silicon substrates has been used to determine the phonon transmission across the metal–silicon interface. A simple rate equation model is used to fit the electron temperature as a function of applied field. From the fit we have established that the phonon transmission deviates from the usual acoustic mismatch model. We find the phonon escape to comprise two components: ballistic escape of phonons from the film and strong scattering close to the interface.

Recent studies of electron heating effects in thin metallic films have been used to measure the low temperature inelastic lifetime for the electrons [1, 2]. In these experiments a small electric field applied across the sample causes the electron system to heat above the ambient temperature of the phonon bath. The inelastic lifetime is then determined from the power dissipated in the film, the electron temperature and the specific heat of the electron system [3]. A necessary assumption when ascertaining the inelastic lifetime in this way is that the phonon system in the film remains in thermal equilibrium with the substrate. Another way to express this assumption is that phonons are expected to escape from the film on a time scale much less than electron–phonon scattering time. It is more instructive to consider the true situation as a bottleneck in which the phonon escape time, t_{esc}, and the electron–phonon interaction time, τ_{ep}, act in series; whichever is the slowest will dominate the relaxation of energy from the electron system to the substrate.

In this letter we describe an electron heating experiment in which the phonon escape time to the substrate is greater than the electron–phonon relaxation time. By using an elementary model to provide a fit to the data, we are able to estimate the magnitude and frequency dependence of the phonon transmission across the film–substrate interface.

The samples consisted of 30 nm thick $Au_{60}Pd_{40}$ alloy films that were thermally evaporated onto an optically smooth silicon substrate at room temperature and at a pressure of 5×10^{-6} Torr. Prior to film deposition, no attempt was made to remove the SiO_2 layer that is known to exist at a silicon surface. The sample geometry was defined by electron beam lithography [4] and had a four-terminal structure with an active region 1 μm wide and 10 μm long. The experiments were performed with the samples under vacuum and thermally anchored to a copper block attached to a temperature controlled bath of liquid helium. Four-terminal resistance measurements of the samples were made over the temperature range 1–4 K using a low frequency AC bridge [5] with a RMS sensing current $\simeq 3$ μA. The fractional resistance change detectable with this arrangement was 2 parts in 10^5. To provide electron heating, a DC field of up to 5 V cm^{-1} was applied

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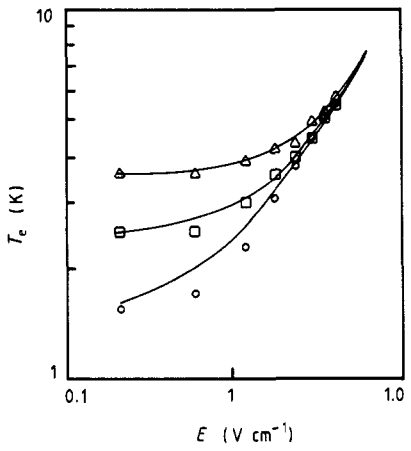

Figure 1. The electron temperature as a function of the applied electric field, E, for a 30 nm thick Au–Pd film on silicon. The substrate temperature was constant for each set of data: 1.5 K (○), 2.5 K (□) and 3.5 K (△). Full curves show the fits to the data obtained using the Pippard electron–phonon interaction and a frequency dependent phonon escape time.

across the samples. The AC sensing current was always kept small enough that its contribution to the heating could be neglected.

Below 4 K we observed an increase of the resistance with a temperature dependence similar to that reported by other workers [6]. From magnetoresistance and Hall measurements of similar Au–Pd films, McGinnis and Chaikin [7] have shown that the low temperature resistance rise can be attributed mainly to weak localisation. To use the temperature dependence of weak localisation as a thermometer for the electron temperature, T_e, the electron–electron relaxation time, τ_{ee}, must be shorter than the electron–phonon relaxation time, τ_{ep}. *A priori* we do not know the relative magnitudes of τ_{ee} and τ_{ep}.

Electron heating data for a Au–Pd film, with a 4 K resistivity of 23 $\mu\Omega$ cm, are shown in figure 1. Low values of the electric field cause little heating above the lattice temperature while at fields above 1 V cm^{-1} the electron temperature becomes independent of the initial lattice temperature. For comparison, in figure 2 we show some previously published data for a similar alloy film deposited on Si_3N_4 [2]. The low electric field behaviour for the two films is the same while the high field behaviour approaches two distinctly different asymptotes.

The asymptotic behaviour of the heating experiments is governed by the bottleneck between the phonon escape time and the electron–phonon inelastic relaxation times. To describe this behaviour we have used a model that was first applied to phonon generation in heat pulse experiments by Perrin and Budd [8]. In the steady state, the same model directly applies to our DC heated metal films. The model is based on the rate equation

$$\frac{dn(\omega)}{dt} = \frac{n(\omega, T_e) - n(\omega)}{\tau_{ep}} + \frac{n(\omega, T_0) - n(\omega)}{t_{esc}} \tag{1}$$

where dn/dt is the net rate of change of the non-equilibrium phonon occupation $n(\omega)$. The Bose distributions of the phonons at the electron temperature and the substrate temperature, T_0, are $n(\omega, T_e)$ and $n(\omega, T_0)$ respectively. Quite simply, the model describes the sum of the net rate of phonons being generated by the heated electrons and the net rate of the subsequent phonon loss to the substrate. In steady state, the sum of the two rates is zero, which allows the non-equilibrium phonon distribution $n(\omega)$ to

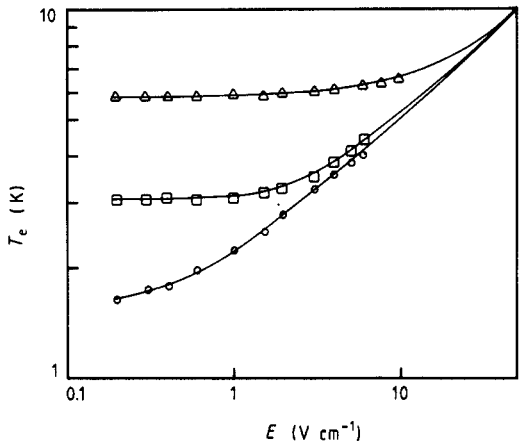

Figure 2. Data taken from [2] showing the electron temperature as a function of the applied electric field, E, for a 30 nm thick Au–Pd film on Si_3N_4. The substrate temperatures were 1.6 K (○), 3.4 K (□) and 6.0 K (△). The full curves are the fits to the data using the present model with the Pippard electron–phonon coupling and a frequency independent phonon escape time.

be expressed in terms of τ_{ep} and t_{esc}. Equating the input electrical power per volume to the power per volume being transferred from the electrons to the phonons, one obtains

$$E^2/\rho = \int \frac{d\omega\, D(\omega)\, \hbar\omega(n(\omega, T_e) - n(\omega, T_0))}{(\tau_{ep} + t_{esc})} \tag{2}$$

where E is the electric field across the metal film, ρ is the film resistivity, and $D(\omega)$ is the phonon density of states. In the present work we have used a Debye density of states for three-dimensional phonons. We note that for thinner films or lower temperatures, the density of states should be modified to account for reduced phonon dimensionality. Since the data to be fitted consist of electron temperatures versus the applied DC electric field at different substrate temperatures, equation (2) can be integrated to obtain a fit once expressions for τ_{ep} and t_{esc} are supplied.

The films used in this and the previous work are of relatively low disorder. From the Drude expression for the conductivity we estimate the elastic mean free path l for the present sample to be 4 nm. Using the dominant phonon wave vector, $q = 2kT/\hbar v_s$, we find $ql \sim$ unity for the temperatures used in this work, and so we have taken the complete form of the Pippard model [9] to describe the electron–phonon interaction

$$\tau_{ep}^{-1} = \frac{Nmv_F}{\rho_D l}\left(\frac{1}{3}\frac{q^2l^2 \tan^{-1} ql}{ql - \tan^{-1} ql} - 1\right). \tag{3}$$

For Au–Pd we have used the electron rest mass for m, a mass density $\rho_D = 16\,400$ kg m^{-3}, an electron density $N = 5.8 \times 10^{28}$ m^{-3}, a Fermi velocity $v_F = 1.4 \times 10^6$ m s^{-1} and a sound velocity $v_s = 1.8 \times 10^3$ m s^{-1}. Using the model of Eiler [10] for the electron–electron phase breaking time we estimate $\tau_{ee} \simeq 4 \times 10^{-10}$ s at 1 K. Comparing the magnitudes of τ_{ee} and τ_{ep} in a dominant phonon approximation, we find $\tau_{ep} > \tau_{ee}$ for the temperatures used in the experiments. Therefore, we are assured that the equilibrium temperature profile of the sample resistance provides us with a well calibrated thermometer of the electron temperature.

Phonon escape from the film has been described by the usual frequency independent escape time

$$t_{esc} = 4\eta d/v_s \tag{4}$$

where d is the film thickness and $\eta = 1/2\Gamma$ with Γ the acoustic mismatch parameter calculated by Little [11].

Substituting these relaxation times in equation (2) and performing the integration, we obtain a fit to the data for the film supported on the Si_3N_4, as shown in figure 2. A value of $\eta \simeq 1$ is used in the fit and we note this is smaller than predicted by Little. Similar discrepancies have been observed by Marx and Eisenmenger [12].

The present data (figure 1) cannot be modelled with just a frequency independent phonon escape time. At low frequencies we find the energy relaxation bottleneck is dominated by the electron–phonon interaction while at high frequencies the phonon escape time dominates and becomes much greater than predicted by the acoustic mismatch model: that is, high frequency phonons appear to be trapped within the film. If the trapping mechanism were intrinsic to the film we would expect to observe the frequency dependent escape time in both the present and the previous work since the films were electrically very similar. We conclude, therefore, that the acoustic properties of the film–substrate interface determine the frequency dependence of the phonon escape time. This opinion is supported by the work of Dolan and Osheroff [6] who mentioned the fact that in their experiments electron heating showed a dependence on the substrate material.

To fit the data we have invoked a frequency dependent acoustic mismatch parameter

$$\eta(\omega) = \eta_0 + \eta_\omega. \tag{5}$$

We use a power law, $\eta_\omega = c\omega^D$, to describe the frequency of the acoustic mismatch parameter and find from the best fit to the data that $c \simeq 8 \times 10^{-50}\ s^4$ and $D \simeq 4$.

The frequency dependent phonon escape time causes the phonon distribution within the metal film to be non-thermal. Using the electron–phonon scattering rate given by equation (3) and the frequency dependent escape time, we have calculated the phonon distribution within the film shown in figure 3. The pronounced peak at 2 THz, corresponding to a phonon wavelength $\simeq 1$ nm, demonstrates that phonons with frequencies above $\simeq 1$ THz are trapped in the film for times longer than calculated from acoustic mismatch theory. This observation is in agreement with pulse measurements of the heat flow across a gold–SiO_2 interface which showed a larger than expected thermal boundary resistance for phonons in the 2–3 THz range [13]. Enhanced trapping of high frequency phonons is also consistent with measurements of the spectrum of phonons emitted from a heated metal film [14].

The physical origin of the frequency dependence is as yet unknown. However, we conjecture that one possible source would be the scattering of short wavelength phonons in the SiO_2 beneath the metal film. Phonons with wavelengths close to those at the distribution maximum have a scattering rate in glasses that is proportional to ω^4 [15, 16]. If we assume the frequency dependent escape time to be the time of phonon diffusion across the oxide layer, this mechanism would provide the frequency dependence we observe. Further experiments are under way to identify the source of the frequency dependence.

In summary, we have presented electron heating experiments in which the steady state temperature of the electron system is dominated by the rate of phonon escape to the substrate. Using a simple model we have obtained a fit to the data which indicates

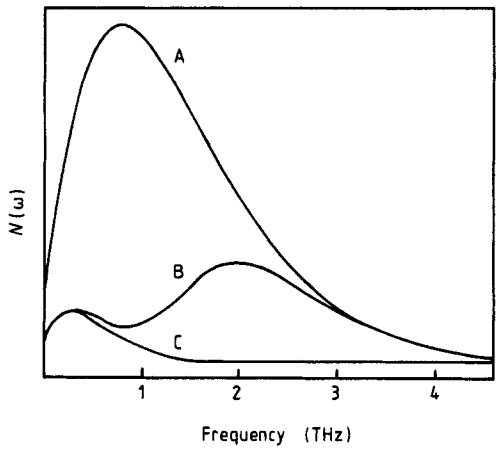

Figure 3. The phonon distributions $N(\omega) = D(\omega)n(\omega)$ as a function of the phonon frequency for a substrate temperature of 1.5 K and an electron temperature of 4.0 K (an applied field of $\simeq 1.8$ V cm^{-1}). For each curve $D(\omega)$ is the Debye density of states and $n(\omega)$ is: (A) the Bose–Einstein distribution at the electron temperature; (B) the calculated non-equilibrium phonon distribution in the film; (C) the Bose–Einstein distribution at the lattice temperature.

that a frequency dependent phonon scattering mechanism exists at the metal–silicon interface. Finally, these data show that attention must be given to the substrate material when performing heating experiments on small structures.

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