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STRIPS AS POSITION SENSITIVE PARTICLE
DETECTORS

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Abstract

The feasibility of superconducting strips as position-sensitive detectors for α -particles was investigated. For this purpose we have prepared films of different material and dimensions using standard evaporation and laser ablation in a UHV system. Films were patterned into strips and current-biased strips were irradiated with 5.5MeV α -particles. The energy deposited by the particle in the strip creates a normal-conducting region, which in turn causes a voltage drop. The time evolution of the voltage drop across the strip as a function of bath temperature and bias current has been investigated. Information about the site of the strip hit by the α -particle was obtained from monitoring the propagating normal-conducting zones. A lateral position sensitivity of $\pm 25\mu\text{m}$ for a Tantalum strip detector has been achieved.

A model describing the time evolution of a normal-conducting zone in a current-biased superconducting strip was used to test and interpret our experimental results. The comparison between model calculation and experiment allows to evaluate the thermal conductivity of the strip material and the heat transfer from the strip into the substrate.

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1. Introduction

The use of superconducting strips for the detection of nuclear particles was already proposed in 1962 by Sherman [1]. More recently superconducting strip detectors have been considered for the realisation of radiation-hard vertex detectors for applications in future high-luminosity colliders. A suitable detection of 5.5MeV- α -particles has previously been demonstrated with In [2,3,4], Sn [2,3,5,6], and Al [7] strips. Recently Gabutti et al. demonstrated the feasibility of granular Nb [8], Al [9], and W [10,11] strips, for monitoring 6keV X-rays. Gray et al. [12] suggested that superconducting strips can provide information on the site of the irradiation impact. Because this assumption has never been verified experimentally, we started to investigate the position sensitivity of Sn and Ta strips for α -particles and X-rays.

The energy deposited in a superconducting strip hit by an α -particle or radiation causes, at the site of the impact, a local temperature rise with small spatial extension (Figure 1). Depending on the amount of deposited energy, the bath temperature at which the strip is kept in equilibrium, and the specific heat of the strip material, the local temperature will eventually rise above the critical temperature T_c , resulting in a normal-state hot-spot (HS). If the size of this initial HS extends over the whole width of the strip, any given bias current I_{bias} produces a voltage drop across the normal-region.

If the size of the initial HS is smaller than the width of the strip, an applied bias super-current is forced to flow in a reduced cross-section of the remaining superconducting material, thereby increasing the current-density. Depending on bias current and HS-size, the critical current-density may be exceeded and consequently the HS spreads over the full strip width, again resulting in a voltage drop, that can be measured.

Once a normal-region (NR) is created over the full width of the strip, this NR may spread along the strip (turning the whole strip length normal) or collapse again, depending on whether the Joule heat produced by the bias current is larger or smaller than the heat transferred per time interval to the surroundings (substrate, exchange gas, etc.). Thus, the time evolution of a NR in a current-biased superconducting strip depends on the thermal and transport properties of the strip material and thermal coupling of the strip to its environment. In this work we only report on NR's propagating over the whole length of the strip. From the propagation time, information can be obtained about the position on the strip where the radiation-induced energy deposition occurred. The best spatial resolution for identifying the site of radiation impact we obtained so far with a Ta strip was 25 μm .

In previous theoretical models on NR-propagation it was assumed that the thermal properties of the strip material like the specific heat c_p and the heat conductivity κ are temperature independent and equal in the normal- and superconducting state [13]. This had the advantage of yielding an analytical solution of the heat conduction equation. In order to treat the problem more realistically, our analysis takes fully into account the temperature dependence of $c_p(T)$ and $\kappa(T)$, and we solve the equations numerically. As will be shown below, good agreement between the experimental NR-propagation

Figure 1: Schematical principle of operation of a superconducting strip detector.

times and the calculated values can be obtained in this way for Sn and Ta strips irradiated with 5.5MeV α -particles.

In the following section we will explain our model for calculating the time evolution of propagating normal-regions. Section 3 describes the sample preparation and the experimental set-up. In section 4, the experimental results are presented, and the numerical solutions of the heat-conduction equation in comparison with experimental data are discussed in section 5.

2. Time Evolution of a Normal-Conducting Zone in a Superconducting Strip (Theory)

The propagation of a normal-conducting zone in a current-biased superconducting strip is governed by the one-dimensional heat-conduction equation

$$\frac{\partial}{\partial x} \left(\kappa(T) \frac{\partial T}{\partial x} \right) - c(T) \frac{\partial T}{\partial t} + p_{\text{in}} - p_{\text{out}} = 0, \quad (1)$$

with

- x : position along the strip,
- t : time,
- $\kappa(T)$: thermal conductivity of the strip material,
- $c(T)$: specific heat of the strip material per unit volume,
- $T=T(x,t)$: position and time dependent strip temperature,
- p_{in} : input-power density per unit volume of the strip,
- p_{out} : power-loss density per unit volume of the strip.

In Eq. (1) we assumed that the strip temperature is constant over the cross-section of the strip and depends only on the position along the strip. This is justified by the small dimensions of our strips in y (width) and z (thickness) directions, respectively.

Broom and Rhoderick [13] previously investigated the time evolution of a normal-to-superconducting boundary (NS-boundary). They considered an infinitely long current-biased strip with, initially, one half being normal-conducting (N), and the other half superconducting (S) (BR-model). The input-power density is equal to the power dissipated by the bias current I_{bias} in the N-region per unit volume. In the N-half of the strip and far away from the NS-boundary the strip temperature approaches the steady state value, given by

$$T_{\infty} = T_{\text{bath}} + \frac{I_{\text{bias}}^2 r}{\alpha w}, \quad (2)$$

where T_{bath} denotes the bath temperature, r is the normal-conducting resistance per unit length of the strip, α is the heat transfer coefficient from the strip to the substrate and w is the strip width. The power-loss density p_{out} is assumed to depend in the following way on the strip and bath temperature

$$p_{\text{out}} = \frac{\alpha(T - T_{\text{bath}})}{d}, \quad (3)$$

where d is the strip thickness. In order to solve Eq. (1) analytically, Broom and Rhoderick [13] assumed that the specific heat and the thermal conductivity are temperature independent and have

the same value in the normal- and superconducting state. From these approximations an analytical expression for the velocity v_{NS} of the NS-boundary has been obtained:

$$v_{NS} = \frac{\sqrt{\alpha \kappa / d}}{c} \times \frac{(T_{\infty} - 2T_c + T_{\text{bath}})}{\sqrt{(T_c - T_{\text{bath}})(T_{\infty} - T_c)}}. \quad (4)$$

In their approximation for the propagation time of an NR in Lead strips, Broom and Rhoderick [13] used the normal-state values for c and κ at the critical temperature. The results obtained by solving Eq. (4) were of the same order of magnitude as experimental values [14]. This procedure may be adequate in certain cases, if only a crude estimate is needed. However, if more accurate values are required, Eq. (1) must be solved more precisely.

Since κ and c are temperature dependent with different functional dependencies in the normal- and superconducting state, respectively, Eq. (1) can no longer be solved analytically. In order to treat the problem more realistically, the following analysis takes fully into account the temperature dependence of c and κ . This can only be done numerically. In the absence of external magnetic fields and assuming that the thermodynamic critical field B_c takes the form [15]

$$B_c(T) = B_c(0) \times \left[1 - \left(\frac{T}{T_c} \right)^2 \right], \quad (5)$$

and the specific heat $c_s(T)$ per unit volume of a superconductor as function of temperature can be written as [15]:

$$c_s(T) = c_N(T) + \frac{2 B_c^2(0) T}{\mu_0 T_c^4} (3T^2 - T_c^2), \quad (6)$$

where $c_N(T)$ is the specific heat in the normal-state, $B_c(0)$ is the thermodynamic critical field at zero temperature, and $\mu_0 = 4\pi \times 10^{-7} \text{Vs/Am}$ is the vacuum magnetic permeability. The specific heat in the normal-state $c_N(T)$ at low temperatures may be approximated by the sum of an electronic and a lattice contribution and takes the form [16]:

$$c_N(T) = \gamma T + \beta T^3. \quad (7)$$

The values of γ and β for Sn and Ta are given in Table 1.

	T_{c0} [K]	Δ_0 [meV]	B_{c0} [mT]	$\gamma \left[\frac{\text{J}}{\text{m}^3 \text{K}^2} \right]$	$\beta \left[\frac{\text{J}}{\text{m}^3 \text{K}^4} \right]$
Sn	3.72	0.58	30.9	109.3	16.12
Ta	4.47	0.70	83	543.9	12.2

Table 1: Superconducting and thermal properties of Sn and Ta taken from ref. [16, 30]

Another important quantity entering Eq. (1) is the thermal conductivity. For a clean normal-conducting metal in the temperature range of interest the heat transport is dominated by the motion of electrons and the lattice contribution can be neglected. The heat-resistivity in the normal-state can then be written as [17]:

$$\kappa_N^{-1} = AT^2 + \frac{B}{T}. \quad (8)$$

Here A and B determine the heat-resistivity due to scattering of electrons by phonons and by lattice defects, respectively. The constant A is an intrinsic property of the metal and can be taken from the literature [18]. The value of B may be evaluated from the measured electrical resistivity ρ using the Wiedemann-Franz-law:

$$B = \frac{\rho}{L_0}, \quad (9)$$

with the ideal Lorenz number $L_0 = 2.45 \times 10^{-8} \text{W}\Omega/\text{K}^2$ [17].

Equation 9 assumes that the electronic mean-free-path values for the electric (l_{el}) and thermal (l_{th}) motion are equal. However, a more accurate solution of the Boltzmann transport equation in the presence of an electric field and a temperature gradient in the sample shows that the values of l_{el} and l_{th} are different [19]. In pure metals and at low temperatures l_{th} may be much shorter than l_{el} . The renormalized Wiedemann-Franz law then reads as [19]:

$$\frac{\kappa_N}{\sigma T} \equiv \frac{\kappa_N \rho}{T} = L = L_0 \frac{l_{th}}{l_{el}}. \quad (10)$$

The normalized value of L has to be inserted for L_0 in Eq.(9). As the ratio $\Phi = l_{th}/l_{el}$ is not known for our strips we used the factor Φ as a fit parameter in our calculations.

In a weakly-coupled superconductor the electronic contribution to the heat conduction is due to quasi-particles, the number of which decreases exponentially with decreasing temperature [15]. The Cooper-pairs by definition do not contribute to the heat transport. The ratio between the electronic thermal conductivity in the superconducting and the normal-state, respectively, can then be derived from ref. [20]:

$$\frac{\kappa_s}{\kappa_N} = \frac{2 F_1(-\zeta) + 2 \zeta \ln(1 + e^{-\zeta}) + \zeta^2 / (1 + e^{\zeta})}{2 F_1(0)}, \quad (11)$$

with $\zeta = \Delta_0 / k_B T$
and $F_1(-\zeta) = \int_0^\infty \frac{\eta d\eta}{(1 + e^{\eta + \zeta})},$

where Δ_0 is the superconducting gap parameter at $T=0K$ and k_B is the Boltzmann constant.

The main heat source in our strips is the power dissipated by the bias current I_{bias} in the normal-region of the strip. This can be modelled by introducing a Heaviside step-function:

$$\Theta(x) = \begin{cases} 0 & \text{for } x \leq 0 \\ 1 & \text{for } x > 0 \end{cases} \quad (12)$$

and describing the power-dissipation density by

$$p_{\text{in}} = \frac{I_{\text{bias}}^2 r(T)}{w d} \Theta(T - T_c(I_{\text{bias}})). \quad (13)$$

Here we neglect the power dissipation at the NS-boundary itself. Stuiyinga *et al.* [21,22] showed that an additional power dissipation occurs in a region of length Λ_{Q^*} at the NS-boundary. In absence of magnetic impurities in the strip, Λ_{Q^*} can be evaluated from

$$\Lambda_{Q^*} = \left(\frac{v_F l_{\text{el}}}{3} \tau_{Q^*} \right)^{1/2} \quad (14)$$

with $\tau_{Q^*} = \frac{4 k_B T_c}{\pi \Delta_0} \tau_E,$

where v_F is the Fermi velocity, and τ_E the inelastic electron scattering time. Using typical values for Sn ($v_F=1.9 \times 10^8 \text{cm/s}$, $l_{\text{el}}=50 \text{nm}$, $\Delta_0 / k_B T_c = 1.76$, $\tau_E=0.27 \text{ns}$ [21,22]), Eq. (14) yields $\Lambda_{Q^*} = 2.5 \mu\text{m}$. Analogous data for Ta are not available, but Λ_{Q^*} is expected to be shorter than in Sn [21,22]. The typical length of our strips ranged from $500 \mu\text{m}$ to 1.5mm . Therefore the apparent prolongation of the normal-region is less than 1% of the total strip length and can be neglected.

The power-loss density may have three different contributions:

- heat transfer due to convection in the surrounding exchange gas
- heat transfer to the substrate
- heat radiation

The first contribution can easily be neglected since we operated our strips in vacuum ($<10^{-3} \text{mbar}$).

	ref. [24]	ref. [25]	ref. [26]
Sn/SiO ₂	28	32	28
Ta/Al ₂ O ₃	---	15	---

Table 2: Heat-transfer coefficient h_0 in $\frac{\text{mW}}{\text{cm}^2\text{K}^4}$ taken from ref. [24, 25, 26].

The heat transfer $\delta Q/\delta t$ across the interface of two solids of temperature T_1 and T_2 can be calculated with the relation [23]

$$\frac{\delta Q}{\delta t} = F h_0 (T_1^4 - T_2^4), \quad (15)$$

in which F is the interface area and h_0 is the so-called heat-transfer coefficient. Theoretically h_0 can be calculated assuming a perfect interface between the two materials using the acoustic mismatch model (AM-model). In this case the phonon transfer rate depends only upon the acoustical and the mechanical properties of the materials in contact (density, density of states and group velocity of the acoustic phonons). In Table 2 we give different values for h_0 for Sn/SiO₂ and Ta/Al₂O₃ interfaces taken from the literature which were calculated by means of the AM-model [24,25,26]. A realistic interface between two solids, however, can be very different from an ideal one. In that case the heat-transfer coefficient is lower, as was experimentally established by Cheeke et al. [24]. Thus the values in Table 2 are an upper limit for h_0 .

The third contribution to the power-loss density is the power radiated off the heated strip which can be neglected due to the very small value of the Stefan-Boltzmann constant $\sigma_{\text{SB}} = 5.67 \times 10^{-9} \frac{\text{mW}}{\text{cm}^2\text{K}^4}$, if compared to the h_0 values listed in Table 2.

The heat-conduction equation, Eq. (1), can thus be rewritten as:

$$\frac{\partial}{\partial x} \left[\kappa(T) \frac{\partial T}{\partial x} \right] - c(T) \frac{\partial T}{\partial t} + \frac{I_{\text{bias}}^2 r(T)}{w d} \Theta(T - T_c(I_{\text{bias}})) - \frac{h_0 (T^4 - T_B^4)}{d} = 0 \quad (16)$$

where the specific heat $c(T)$ from equations 6 and 7 and the thermal conductivity $\kappa(T)$ from equations 8 and 11 have to be inserted.

Simulations of the propagation behaviour of NR in Sn and Ta strips will be presented in section 5 and compared to our experimental data.

3. Samples and Experiment

3.1. Sample Preparation

Tin

The Sn films were deposited on fused quartz substrates in a Balzers UMS630 ultra-high-vacuum (UHV) system using thermal evaporation out of an effusion cell. The background pressure in the UHV-system before starting the evaporation process was typically between 1 and 6×10^{-10} mbar. By regulating the temperature of the effusion cell heater, the deposition rate could be adjusted to between 0.1 and 1 nm/s.

Sn-films grown at pressures below 4×10^{-7} mbar are non-conducting and consist of an assembly of disconnected grains or islands (size $\approx 1 \mu\text{m}$). In order to promote the growth of interconnected islands, oxygen had to be bled into the vacuum chamber during the deposition process. The oxygen pressure

Figure 2: Scanning electron micrographs of two Sn-films:
(a) non-conducting film deposited at 2×10^{-9} mbar
(b) conducting film deposited at 2×10^{-6} mbar O_2

Figure 3: Resistivity at 4.2K and 300K of Sn-films deposited at a rate of 0.1nm/s as function of the oxygen pressure during film growth.

required to produce conducting films ranged from 5×10^{-7} to 1×10^{-5} mbar.

In Figure 2 scanning electron micrographs of two Sn-films are shown. One was deposited at 2×10^{-9} mbar total pressure (a), and one was grown at 2×10^{-6} mbar partial O_2 -pressure (b). The oxygen enhances the sticking probability of the Sn-atoms on the substrate and reduces the island-type growth [27]. It is known that Sn grows with smaller grains on cold substrates ($T \approx 77K$). However, we could not use this technique without damaging the resist stencils prepared on the substrates for the lift-off patterning, because photoresist layers do not withstand extreme high or low temperatures.

Figure 3 shows the resistivity of the Sn-films as a function of oxygen pressure during evaporation. The deposition rate of all films was 0.1nm/s. The films produced at oxygen pressures above 5×10^{-5} mbar consist of interconnected grains but are electrically insulating. This is most likely due to the oxidation of the grain boundaries, or of the whole material forming the grains. The relative resistivity maximum at $8 \cdot 10^{-7}$ mbar is not understood.

The patterning of the Sn-films was done using a photolithographic lift-off process. The strips obtained were typically 0.5 to 1.5mm long, 3 to 12μm wide and 100 to 500nm thick.

Tantalum

The Ta-films were deposited in the same UHV-system as the Sn-films on $\langle 11\bar{2}0 \rangle$ -oriented Al_2O_3 -substrates. For the deposition of the Ta-films we used the pulsed-laser-deposition (PLD) method. The experimental set-up of our laser-ablation system is shown in Figure 4. The beam of a high-power excimer-laser (maximum pulse energy: 0.8J; maximum repetition rate: 10Hz; gas mixture: ArF;

Figure 4: Schematic diagram of the laser-ablation system.

wave-length: 192nm) was focused on a Ta-target placed in the UHV-chamber. In order to avoid significant energy losses we used a UV-grade fused silica lens and a MgF_2 vacuum window. The total laser-target transmission coefficient is 50%. The reflectivity of Ta at 192nm is also 50% [28], thus the total energy deposited on the target is 25% of the initial pulse energy.

The background pressure in the UHV-chamber prior to the evaporation process was again typically between 1 and 6×10^{-10} mbar. Focusing the laser beam on a virgin spot of the target increases the pressure in the evaporation system to about 5×10^{-8} mbar. The main residual gases in the system were

Figure 5: Deposition rate for laser-ablated Ta-films as a function of the laser fluence on the target.

hydrogen and water vapour, as determined by means of a quadrupole-mass-spectrometer. This implies that the surface of the Ta-target is covered by an adsorbed layer of water and hydrogen. This surface layer had to be removed before film growth, in order to minimize contamination of the film. This was achieved by 'cleaning' the target surface with several thousands of laser pulses after each venting of the vacuum system.

Figure 6: Temperatures $T_c^{10\%}$ and $T_c^{90\%}$ and residual-resistivity ratios RRR of Ta-films made by laser-ablation as a function of the energy density E_D of the laser beam on the target. The thicknesses of the films represented in this plot range from 100nm to 300nm.

We operated the laser at its maximum pulse energy of 0.8 J/pulse ($\pm 10\%$) and the repetition rate was varied between 1 and 10 pulses/s. The substrates were held at room temperature during deposition, and placed at a distance varying between 10 and 20 mm vis à vis the target. Due to the narrow material plume induced by the laser radiation the thickness profile of the film is gaussian-like with a FWHM of about 3mm as determined with profilometer measurements.

In Figure 5 we show the dependence of the deposition rate on the energy density E_D of the laser beam on the target. There is an obvious threshold $E_D = 10 \text{ J/pulse} \times \text{cm}^2$ which has to be exceeded in order to initiate the deposition process. Above the threshold, the deposition rate increases with increasing E_D up to 0.2nm/s at 100 J/pulse $\times$ cm². For higher E_D the deposition rate remains constant. Increasing the target-substrate distance by a factor of 2 reduces the deposition rate by a factor of 4 as may be expected from geometrical considerations.

Figure 7: Scanning electron micrograph of a laser-ablated Ta-film.

It has previously been found that in order to initiate the nucleation of solid Ta in the attempted superconducting α -phase by e-gun evaporation on $\langle 11\bar{2}0 \rangle$ -oriented Al_2O_3 -substrates, a minimum substrate temperature of 400°C or a thin Nb underlayer as a crystal growth template is required [29,30]. However, due to the relatively high kinetic energy available in the plasma of the PLD process we obtained Ta films with critical temperatures close to the bulk value of 4.47K on sapphire substrates kept at room temperature without using any growth template.

The resistivity of the films was measured with a 4-point technique at a negligible bias current-density of 1 to $10\text{A}/\text{cm}^2$. The transition width was defined as the difference between the temperatures $T_c^{10\%}$ and $T_c^{90\%}$. $T_c^{10\%}$ and $T_c^{90\%}$ denote here the temperatures at which the resistivity decreased to 10%, respectively to 90% of its value in the normal-state above the onset. The critical temperature T_c was defined as the temperature at which the resistivity decreased to one half of its constant value in the normal-state, just above the onset.

Figure 6 shows the E_D dependence of the critical temperature (upper part) and the residual resistivity ratio RRR (lower part). The RRR increases with increasing E_D to a maximum value of 5 above $200\text{J}/\text{cm}^2$. The critical temperature shows a similar trend: the measured critical temperature approaches the bulk value of 4.47K above a fluence $E_D \approx 200\text{J}/\text{cm}^2$ and the transition width decreases from several degrees at $10\text{J}/\text{cm}^2$ to about 0.1K at higher fluences.

The maximum critical temperature of our films was 4.45K, which is not significantly lower than the bulk value cited in the literature [30]. From the low T_c and the large transition width observed at lower E_D it is quite likely that the films do not consist of one single phase, if prepared under these conditions. This was confirmed by X-ray analysis of the films. Ta-films made at low fluences consist of both α and β -phase. The Ta- β -phase has a T_c of less than 0.5K and its crystal structure is fcc. The α -phase in contrast has a bcc-structure and a T_c of 4.47K. Increasing E_D continuously decreases the amount of β -phase down to less than 1% of the amount of α -phase.

Figure 8: Temperature dependence of the critical current of two different strips (solid squares and circles) and the theoretical temperature dependence as predicted by Skocpol [31] (solid line), and by the Ginsburg-Landau theory [15] (dashed line), respectively.

As can be seen in Figure 7 the surface of a laser-ablated film is not smooth but plagued by sub-micron diameter particles. This is a clear disadvantage when multilayer engineering is envisaged. Another disadvantage of the PLD-process is the thickness inhomogeneity of the films.

The patterning of the Ta-films was done using photolithography and reactive ion etching in an Ar/SF₆ gas-mixture. The strips obtained were typically 0.5 to 1.5mm long, 3 to 12μm wide and 100 to 500nm thick.

3.2. Critical current

Since strip detectors are operated with a bias current, the knowledge of the critical currents of those films is of importance. The critical currents of the superconducting strips were determined by slowly increasing the bias current until a voltage larger than 5μV developed across the strip. The strips used here display type I superconducting behaviour [15]. The voltage rises from below 5μV to the normal-state value within less than 5μA. For comparison the typical values of the critical current in our strips ranged between 10mA and 60mA at 2.2K. The temperature dependence of the measured critical current I_c displays the expected quadratic behaviour for wide strips [31] namely:

$$\frac{I_c}{I_{c0}} = 1 - \left(\frac{T}{T_{c0}} \right)^2. \quad (17)$$

Here I_{c0} is the critical current at zero temperature, and T_{c0} the critical temperature at zero current and zero external field. In Figure 8 the critical current of an 11 μm wide and 220nm thick Sn strip and of a 3 μm wide and 340nm thick Ta strip are shown as a function of temperature. The observed temperature dependencies indicate that the current-density in the strips is peaking at the edges as was shown by Skocpol [31]. A homogeneous current-density distribution is expected to result in a critical current proportional to $(1-t^2)^{3/2}(1+t^2)^{1/2}$ with $t = T/T_{c0}$ [15], as indicated by the broken line in Figure 8.

The maximum current that can be passed through the strips before they switch to the normal-state is limited by the segment of the strip having the smallest cross-section. Effects due to such inhomogeneities are expected to be more pronounced for Sn granular films than for Ta layers, which exhibit a more homogeneous morphology. However, inhomogeneous etching, due to the thickness inhomogeneity of the Ta films produced by laser-ablation may also give rise to a cross-section variation of $\pm 5\%$ along the strips.

3.3. Experimental Set-up

For the radiation-monitoring experiments, the strips were irradiated with 5.5MeV α -particles emitted by a ^{241}Am -source. The amount of energy deposited in the strips depends, of course, on the material and can be estimated using [32]:

$$\begin{aligned} \frac{dE}{dx}(\text{Sn}) &= 150 \frac{\text{keV}}{\mu\text{m}} \\ \frac{dE}{dx}(\text{Ta}) &= 340 \frac{\text{keV}}{\mu\text{m}} \end{aligned} \quad (18)$$

The source and the strip were mounted opposite to each other and 2mm apart in a vacuum container. Since the active area of the source was 4 mm in diameter, the irradiation of a 12 μm wide and 1.5mm long strip can safely be assumed to be uniform. The container was immersed into a ^4He bath, the temperature of which could be varied between 2.17 and 4.2K. A fast GaAs-preamplifier, held at bath temperature, was used to match the strip impedance to a 50 Ω coaxial line, connecting to the room-temperature electronics. As current-source we used a GaAs-FET which was mounted on the same chip as the preamplifier. The maximum current provided by this FET-current-source is 30mA.

Figure 9: Block diagram of the read-out electronics to measure α -particle-induced voltage signals in current-biased superconducting strips.

Figure 9 shows the schematic diagram of the read-out electronics. The voltage signal was displayed on a 350MHz digital oscilloscope and the rise time was read out by a computer connected to the oscilloscope. The rise time of the voltage signals was obtained by measuring the time lapse between the 10% and the 90% magnitude transition points. A discriminator was used to trigger the electronics, as the voltage signal exceeded a threshold value of 50mV.

Since, by definition, a spreading NR drives the full length of the strip normal, the bias current has to be interrupted to allow the strip to regain full superconductivity and to cool again to bath temperature. Therefore the current-source was switched off 20 μ s after the detection of a voltage signal. This time was long enough to allow the NR to extend over the full length of the strip (NR spreading times were typically shorter than 10 μ s) but also short enough to avoid overheating of the strip. The cooling time (i.e. the time during which the current-source was switched off) was typically 1 ms for the largest bias currents used (30mA) and could be reduced to approximately 100 μ s for smaller currents ($\approx$ 1mA). The count rate that will be presented below is dead-time corrected.

4. Experimental Results

4.1. Radiation Detection

If an X-ray is absorbed in the strip its full energy of 5.9keV is deposited. All X-rays which pass the strip are absorbed in the substrate. The amount of energy transferred from the substrate into the strip by phonons can be neglected. In the case of an α -particle impact there are two different contributions to the energy available in the strip for the hot-spot production. The first contribution is the energy directly deposited by the α -particle passing through the strip. It may be calculated using Eq.(18) and the thickness of the strip. After passing the strip, the α -particle enters the substrate where it is stopped after a few μm . One part of the energy deposited in the substrate may reach the strip by phonon transfer and increase the total energy available. The portion of energy transferred from the substrate to the strip may be as much as 5% of the total α -particle energy [33]. Due to the higher energy available in the strip after an α -particle impact, as compared to the absorption of a 5.9keV-X-ray, a wider strip is quenched or, in other words, a lower bias current is needed to quench a strip of a given width. Because of the lower specific heat and the lower critical temperature of Sn the same energy E is able to switch a larger volume into the normal-state than in the case of Ta. Therefore a lower bias current is required to quench the remaining cross-section of a Sn strip than that of a Ta sample of the same width (see section 5). We were not able to detect 5.9keV-X-rays in our wide strips (width $> 3\mu\text{m}$). For a successful detection of keV-X-ray radiation by this technique, narrower strips (width $< 1\mu\text{m}$) are required.

The normal-conducting zone which develops over the full width of the strip can either spread along the strip or contract again and collapse, depending on the balance between the Joule heat produced by the bias current in the normal-region and the heat transfer to the environment. An expanding hot-spot causes an increase of the strip resistance and therefore an increase of the voltage drop along the strip. From the time evolution of the voltage signal, information on the dynamics of the propagating hot-spot can be obtained. Qualitatively, we expect a voltage rise along the strip up to the value corresponding to the normal-state resistance of the whole strip.

In Figure 10 three different α -particle-induced voltage traces are plotted. Curve 2 represents the general case. The α -particle impact occurs somewhere between the middle of the strip and one of its ends. The hot-spot initiated in this way begins to expand along the strip, i.e. the two NS-boundaries, which enclose the normal-region, move in opposite directions, each of them with a velocity v_{NS} . Thus, the length of the normal-region which is directly proportional to the voltage drop along the strip, increases with the velocity $2v_{\text{NS}}$. This yields the first slope in curve 2. After a time interval τ_1 one of the NS-boundaries reaches the nearest contact pad and stops moving further. This is because the contact pads act as heat sinks held at the constant temperature T_B . Due to their dimensions (10mm^2), the Joule heat produced in them is 10^3 times smaller than in the strip. After the time τ_1 , only one moving NS-boundary remains, i.e., the spreading velocity of the normal-region is now only

Figure 10: Oscilloscope traces showing typical α -particle-induced voltage signals across a $5\mu\text{m}$ wide, 200nm thick, and 1.3mm long Sn strip. The fastest signal (curve 1) arises from a hot-spot created in the middle of the strip. Signal 3 is due to an α -particle hitting the strip near one of its ends. In signal 2 the change of slope corresponds to one flank of the hot-spot reaching the end of the strip. The bath temperature was 2.6K and the bias current was 15mA , corresponding to 96% of the critical current at this temperature.

v_{NS} . Accordingly the slope of the voltage signal changes, as may be seen in curve 2 of Figure 10. The remaining NS-boundary reaches the other contact pad at the time τ_r . Thus the full length of the strip is now in the normal-state. After τ_r a slow voltage rise remains, as can be observed in Figure 10. This is due to further heating of the sample.

The other two signals in Figure 10 can now be interpreted straightforwardly. Signal 1 arises from an α -particle impacting exactly in the middle of the strip. The two NS-boundaries reach the contact pads at the same time τ_{min} . The voltage trace has a constant slope which is necessarily parallel to the initial slope of signal 2. Signal 3 is due to a hot-spot initiated very close to one contact pad. The hot-spot expands effectively only in one direction. The slope of the voltage signal equals the slope of signal 2 after time τ_1 , and its rise time τ_{max} is twice the rise time of signal 1. From the constant slopes in Figure 10 it is concluded that the propagation velocity of the NS-boundary is constant.

Figure 11: Rise time of voltage traces for α -particle impacts in the middle of an $11\mu\text{m}$ wide, 220nm thick, and 1.3mm long Sn strip (a) and of a $3\mu\text{m}$ wide, 340nm thick and $650\mu\text{m}$ long Ta strip (b), plotted versus bias current for several bath temperatures. The solid lines are guides to the eye.

We investigated τ_r of particular signals as a function of bath temperature and bias current. In Figure 11 the measured rise times τ_r for events induced by α -particles impinging in the middle of an $11\mu\text{m}$ wide, 220nm thick, and 1.3mm long Sn strip and of an $3\mu\text{m}$ wide, 340nm thick, and $650\mu\text{m}$ long Ta-strip, respectively are plotted versus bias current for several bath temperatures. The rise time increases both with decreasing bias current at constant temperature due to less Joule heat produced in the strip and with decreasing temperature at constant bias current due to more effective cooling of the strip. In the next section numerical calculations of the NR-propagation time will be presented and compared with the experimental data.

Figure 12: Measured count rates in an 11 μm wide, 220nm thick and 1.3mm long Sn strip as a function of bias current for three different bath temperatures. The lines are guides to the eye.

The count rate of the Sn strip detector is plotted in Figure 12 as a function of bias current for different bath temperatures. According to previous measurements in other materials [10], a sharp rise in the count rate at a certain current is expected, followed by a current independent plateau (see inset in Figure 12) implying that the entire strip has become sensitive. However, our experimental results show a broad rise and no plateau. This is primarily caused by the inhomogeneous current-density distribution over the cross-section of the strip. As the bias current is raised from zero, the strip edges first become sensitive to α -particles. A normal-conducting zone forces the bias current to flow through the remaining superconducting cross-section. A normal-zone at the strip edge redistributes a larger amount of current than the same zone in the middle of the strip. Therefore the increase of current-density in the remaining superconducting part of the strip is larger in the first case than in the second. Increasing the bias current enhances the current-density as well and therefore the sensitive part of the strip width. Using the relation for the current-density distribution in an 11 μm wide strip [31] it is estimated that for bias currents larger than 80% of the critical current the full strip width becomes sensitive to α -particles resulting in a current independent count rate at higher currents. The granular structure of our Sn films, however, inhibits the highest current densities to be reached.

The inhomogeneity of the current-density distribution in narrower strips is less pronounced [31]. Hence, the maximum sensitivity can be reached at I/I_c ratios smaller than for wide strips. In narrow strips, however, the inhomogeneity of the cross-section due to the granular structure of the Sn films plays a more important role. The weakest section of the strip may have a cross-section which is up to one order of magnitude smaller than average, causing a very low critical current. In order to achieve a higher sensitivity, strips of more uniform cross-section are required.

Figure 13: Measured count rates in a $3\mu\text{m}$ wide, 340nm thick and $650\mu\text{m}$ long Ta strip as a function of bias current at four different bath temperatures. The lines are guides to the eye.

The cross-sections of the Ta strips were much more homogeneous than those of the Sn strips. This made it possible to fabricate strips narrow enough in order to verify the expectation mentioned in the last paragraph. In Figure 13 we present the count rate of a $3\mu\text{m}$ wide, 340nm thick and $650\mu\text{m}$ long Ta strip as a function of bias current at different temperatures. The edges of the strip which carry the largest current-density are sensitive to α -particles at about 10% of the critical current. Due to the more homogeneous current-density in narrow strips the maximum sensitivity is reached at about 50% of the critical current and the count rate is current independent at higher currents.

We also measured the count rate in a $7\mu\text{m}$ wide Ta strip (Figure 14). The same behaviour as in the wide Sn strips was observed. The efficiency of the detector starts at 20% of the critical current and increases continuously up to I_c . The limited maximum current provided by our current-source made it impossible to measure the count rate up to I_c at the lower temperatures (2.2K and 2.8K). However, there is little doubt that the count rate will display the same behaviour as at 3.5K.

The measured count rates are roughly two times larger than values calculated on the basis of the activity of the source, the distance to the strip and its area. It seems that α -particles, which hit the substrate close to the strip, give rise to hot-spots as well, i.e., phonons produced by α -particles in the substrate diffuse into the strip and drive it normal. In order to check this assumption we calculated the quartz volume, whose temperature would rise from T_B to T_{max} after absorption of 5.5MeV . With $T_B=2\text{K}$, $T_{\text{max}}=10\text{K}$ and considering the specific heat of quartz in the temperature range of interest [34], a quartz volume of $14\mu\text{m}^3$ is obtained which corresponds to a sphere with a diameter of $6.6\mu\text{m}$,

Figure 14: Measured count rates in a 7 μm wide, 200nm thick and 1.5mm long Ta strip as a function of bias current for three different bath temperatures. The lines are guides to the eye.

of the same order of magnitude as the strip width. This may indeed lead to an apparent expansion of the sensitive strip surface as was observed experimentally.

4.2. Position Sensitivity

The dependence of the signal slope (Figure 10) on the site of impact of the α -particle provides a simple way to use these devices as position-sensitive detectors. By measuring the times τ_1 and τ_r of a voltage signal one may obviously evaluate, at which distance x_0 from the middle of the strip the NR was generated. The site coordinate x_0 is given by

$$x_0 = \frac{l}{2} \times \frac{\tau_r - \tau_1}{\tau_r + \tau_1} \quad (19)$$

where l is the length of the strip. In practice the actual position on the strip is determined by measuring the rise time τ_r and using the relation $\tau_{\min} = \frac{1}{2} \times (\tau_r + \tau_1)$. The result reads

$$x_0 = \frac{l}{2} \times \left(\frac{\tau_r}{\tau_{\min}} - 1 \right) \quad (20)$$

where $\tau_{\min}$ is the minimum rise time (of an event which occurred in the middle of the strip) and can be thought of as a calibration parameter of the strip detector, which only depends on material

Figure 15: Voltage-pulse distribution as a function of strip position.

properties and may be measured separately. The values of l and $\tau_{\min}$ are available with very high accuracy. Thus the spatial uncertainty of the strip detector depends only on the experimental uncertainty of τ_r

$$\Delta x_0 = \frac{l}{2\tau_{\min}} \Delta \tau_r = v_{NS} \Delta \tau_r . \quad (21)$$

Since it cannot be distinguished which half of the strip is hit by the particle an actual detector would require the shielding of one half of the length of the strip.

The spatial sensitivity was investigated using a mechanical mask. A 0.2mm thick copper foil containing a 60μm wide slit was placed between the strip and the source at approximately 0.2mm off the middle of the strip and approximately 0.2mm above the strip. Slit and strip were perpendicular to each other. Thus only approximately 60μm of the whole strip length of 1.5mm were irradiated by α-particles, for a chosen geometrical setting.

Figure 15 shows the distribution of events as a function of x_0 which was calculated from the measured rise time spectrum using Eq.(20). The data shown in Figure 15 are measured with a 7μm wide, 200nm thick and 1.5mm long Ta strip operated at 2.8K and a bias current of 30mA. A distinct maximum is observed around 180μm off the strip center. A Gaussian curve with FWHM=57μm, close to the width of the slit, fits these data nicely. An accurate value for the position sensitivity of the strip detector cannot be extracted from these data. For this purpose a two slit mask with slit width less than the expected position sensitivity is required. This experiment is in progress. From the time resolution of our electronics ($\Delta \tau_r \approx 10$ ns) which was mainly due to noise in the electronic circuit and pick-up noise in the cables, an upper limit for the spatial uncertainty of $\pm 25 \mu\text{m}$ can be estimated. Better shielding of the electronics may increase the signal-to-noise ratio, hence improving the spatial resolution to better than $\pm 10 \mu\text{m}$, independent on the strip length. Therefore such strip devices can in principle be used as position-sensitive detectors.

5. Model Calculations

5.1. Production of the Normal-Region

As outlined above the main criterion for the detection of radiation using a current-biased superconducting strip detector is given by the amount of deposited energy required for the production of a large enough HS such that a bias current distinctly lower than the critical current is able to quench the full width of the strip. The dimension of the initial HS depends only on the deposited energy and the specific heat c (Eq. (6)) of the strip material. In order to calculate the minimum current required to switch the full cross-section of a superconducting strip of given width w after deposition of an energy E we used the following assumptions:

- the deposited energy produces a gaussian temperature profile around the site of impact of the particle at the time $t = 0$

$$T(r, t = 0) = T_{\max} \times \exp\left(-\frac{r^2}{\sigma^2}\right) + T_B \quad (23)$$

with r being the distance from the impact site, σ is the width of the temperature profile, $T_{\max}$ is the incremental-temperature amplitude and T_B is the bath temperature.

- Since the width-to-thickness ratio of the strips ranged from 10 to 100 the temperature gradient perpendicular to the strip-substrate-interface can be neglected.
- the current-density is constant over the strip width and the critical current-density j_c may be calculated from the measured critical current I_c^{meas} and the cross-section of the strip $w \times d$ (see also Figure 1) by

$$j_c = \frac{I_c^{\text{meas}}}{wd}. \quad (24)$$

Using these assumptions, the integral of Eq.(22) over the considered volume yields the relation between the deposited energy E and $T_{\max}$

$$E = 2\pi d \int_0^{w/2} r dr \int_{T_B}^{T(r,t=0)} c(T) dT. \quad (25)$$

We assumed here that the volume heated-up by the deposited energy is of cylindrical shape (see Figure 1). Once the spreading normal-region has developed over the whole strip width we assume it to be of rectangular shape (see Figure 1 and next section). Due to the functional dependence of the upper integration limit in Eq.(25) the integral can only be evaluated numerically. Eq.(25) contains two unknown parameters: the maximum temperature increment $T_{\max}$ and the width σ of the

Figure 16: Maximum widths of a Sn and a Ta strip, both with 300nm thickness, which can be quenched by a bias current I_{bias} after absorption of a 5.9keV X-ray and after passing of a 5.5MeV α -particle as a function of the bias current normalised to the critical current I_c . The calculations were done for a bath temperature of 2.2K.

temperature distribution. Since it is not possible to determine both unknown parameters independently we need another assumption in order to approximate the dimension of the initial hot-spot. Therefore we define

$$T_{\text{max}} = 2T_c - T_B. \quad (26)$$

This approximation is only used to yield an estimate for the sensitivity of the strip. As it will be shown later, increasing the amplitude of the initial temperature distribution by a factor of two has a negligible influence on the rise time τ_r . Numerical simulations of the NR-propagation are described in more detail in section 5.2.

The integral of the critical current-density j_c over the strip width yields the current I_Q which is required to quench the full width of the strip

$$I_Q = d \int_{-w/2}^{w/2} j_c(T(y)) dy. \quad (27)$$

If I_Q is larger than the bias current the hot-spot collapses and no voltage signal as shown in Figure 10 can be detected. If the bias current is larger than I_Q the initial hot-spot spreads over the full cross-section of the strip giving rise to a detectable voltage drop across the strip length. The inversion of Eq.(27) yields the maximum strip width $w_{\max}$ which can be driven normal by a given bias current I_{bias} after deposition of an energy E . In Figure 16 we present the calculated bias current dependence of the maximum strip width $w_{\max}$ which is quenched by an α -particle passing through a 300nm thick strip of Sn or Ta, respectively. We also sketch $w_{\max}$ for the absorption of 5.9keV-X-rays.

5.2. Propagation of the Normal-Region

The time evolution of an NR in a current-biased superconducting strip is governed by the heat-conduction equation Eq.(16). In the following the critical current, the Lorenz number, and the heat transfer coefficient are considered as fitting parameters when we compare our data to the model outlined in Section 2.

Since the current-density distribution is peaked at the edges, the critical current of a homogeneous strip of uniform cross-section is smaller than the critical current-density multiplied by the cross-section [31]. In addition, the experimentally determined critical current I_c^{meas} is limited by the segment of the strip with the smallest cross-section. Since for our model calculations we consider only strips of uniform cross-section, the theoretical critical current I_c^{theo} , which is larger than the measured value I_c^{meas} , has to be considered as well. Therefore we have to introduce a factor $i = I_c^{\text{theo}} / I_c^{\text{meas}}$ as an additional fit parameter in our calculations. It may be interpreted as the ratio between the average cross-section and the minimal cross-section.

As outlined before (see Section 1) the value of the Lorenz number in pure metals at low temperatures decreases with decreasing temperature. Therefore, residual resistivity data used to evaluate the electronic contribution to the thermal conductivity yield too large values, at least in the temperature range of interest here. Anticipating a mean value of $L < L_0$ (see Eq.(10)) in our calculations, we introduce a scaling factor Φ ($0 < \Phi < 1$). The thermal conductivity in the normal-state then reads

$$\kappa_N^{-1} = AT^2 + \frac{1}{L_0} \frac{\rho(0)}{\Phi \cdot T}. \quad (28)$$

	i	Φ	$h_0 \left[\frac{\text{mW}}{\text{cm}^2\text{K}^4} \right]$		
	fit	fit	fit	exp. (ref. [35])	AM-model
Sn/SiO ₂	4±0.2	0.8±0.1	15±1.5	12	28
Ta/Al ₂ O ₃	1.7±0.1	0.5±0.1	11±1	---	15

Table 3: Fit parameters of our calculations compared to predictions of the AM-model and to the experimentally evaluated heat-transfer coefficient by Holt [35].

The third fit parameter is the heat transfer coefficient h_0 which is a measure of the thermal coupling between the film and the substrate.

In order to simplify the problem we considered only NR which originate in the middle of the strip. Thus, Eq.(16) has to be solved for one half of the strip only. In this case the boundary conditions take the form

$$\begin{aligned} \frac{\partial T}{\partial x} &= 0 \quad \text{for } x=0 \quad (\text{strip center}) \\ T &= T_B \quad \text{for } x=\frac{l}{2} \quad (\text{strip end}) \end{aligned} \quad (29)$$

In addition to the boundary conditions, we also need to define a starting condition for solving Eq.(16). This enters as the shape of the initial HS and we assume a Gaussian temperature distribution with a full width at half maximum of 3.5μm (Sn strips) or 2.5μm (Ta strips) respectively. These values were obtained by means of the model presented at the beginning of the present chapter (Eq.(25)), assuming a Gaussian temperature profile with a maximum of $2T_c$. Since the amplitude of the initial HS is also an assumption, we estimated its influence on the calculated rise times. Changing the amplitude of the initial temperature distribution from $1.5T_c$ to $4T_c$, causes a change in the rise time smaller than 1%.

The results of the calculations for a 11μm wide, 220nm thick and 1.3mm long Sn strip are shown in Figure 17 together with experimental data obtained at six different temperatures. Fairly good agreement between theory and experiment is obtained. The best fits of Eq.(16) to the experimentally evaluated rise times was obtained for the parameter set shown in Table 3.

Each of the three fit parameters used in our calculations has a different influence on the shape of the theoretical curves in Figure 17. A change of i or Φ shifts the curves along the current or the rise time axis, respectively, while a variation of h_0 results in a different slope of the curves. From this point of view, no other set of parameters (i, Φ, h_0), with physically reasonable values, can fit the experimental data equally well.

Figure 17: Measured rise times of voltage signals for NR's induced in the middle of an $11\mu\text{m}$ wide, 220nm thick, and 1.3mm long Sn strip (full circles). The lines are fits of Eq. (16) to the experimental data.

From our calculations, the critical current of the Sn strip seems to be four times larger than measured. This implies that the weakest section, which is responsible for the measured critical current, is approximately 60nm thick, i.e., four times thinner than the average thickness of the strip. This agrees with SEM pictures of the strips which show coarse-grained structure with up to $1\mu\text{m}$ large grains, rising above a thin Sn layer, the thickness of which is, here and there, less than 100nm .

The value of 0.8 for Φ seems reasonable. Due to the granular structure of the Sn films the defect and surface scattering of the electrons cannot be neglected. Hence the electronic mean free path is reduced and, in the temperature range of interest, the effective Lorenz number of our Sn strips is apparently smaller than the theoretical value. The deviations between the model calculations and the experimental data shown in Figure 17 are most likely due to the fact that we assumed a temperature

Figure 18: Measured rise times of voltage signals for NR's induced in the middle of an 3 μm wide, 340nm thick, and 650 μm long Ta strip (full circles). The lines are fits of Eq. (16) to the experimental data.

independent Lorenz number. Different bath temperatures and bias currents result in a different steady-state temperature of the normal-conducting zone.

The heat-transfer coefficient h_0 extracted from our data, was compared with values published previously in the literature. From the so-called *acoustic mismatch* (AM) model [23], which treats the heat-conduction between two materials via phonons, a value of $28 \frac{\text{mW}}{\text{cm}^2 \text{K}^4}$ is predicted [24,25,26].

Herein, a perfect interface between Sn and quartz is assumed. As shown by Little [23], the heat-transfer coefficient is smaller in the case of an imperfect contact. Cheeke *et al.* [24] measured the power transferred from a Sn film to the quartz substrate as a function of film temperature. In the temperature range of interest their results were roughly a factor of two smaller than predicted by the AM model. Holt [35] measured the thermal resistivity at the interface between Sn and quartz. From these measurements $h_0 = 12 \frac{\text{mW}}{\text{cm}^2 \text{K}^4}$ may be deduced. It is known that the evaporation parameters affect the structure of the films, which in turn affects the matching of the film to the substrate. This was also shown by Cheeke *et al.* [24], who observed a strong dependence of the phonon transfer efficiency upon the substrate temperature during film growth. Taking all this into account, our value for h_0 resulting from the fit procedure is within reasonable limits.

In Figure 18 we show results of rise time experiments and corresponding calculations for a 3 μm wide, 340nm thick and 650 μm long Ta strip. The best fit was obtained with the set of parameters again listed in Table 3. We consider the agreement between experiment and calculation as very good.

As mentioned before the Ta strips have a more uniform and homogeneous cross-section than the Sn strips. Small defects due to non-uniform etching, however, cannot be avoided. Thus the factor $i=1.7$ is reasonable. Ultra-high-vacuum conditions during evaporation are likely to yield Ta films with less

impurities than the Sn films, which are deposited in the presence of O_2 . Therefore the decrease of the Lorenz number in the temperature range of interest can be more pronounced than in the Sn strips. This explains the lower value of Φ in the Ta films.

The heat-transfer coefficient h_0 is smaller than the value of $15 \frac{\text{mW}}{\text{cm}^2 \text{K}^4}$ which is predicted by the AM-model [25], as may be expected in case of a realistic (non-ideal) substrate-film interface. To our knowledge, no other experimental data on the heat-transfer coefficient between Ta and Al_2O_3 are available in the literature.

6. Conclusions

In this work we investigated the physical properties of current-biased superconducting strips and their feasibility as position sensitive α -particle detectors. We focused our attention on the time evolution of normal-conducting regions which expand along the strip. The rise time of the voltage-pulse induced by a spreading normal-region can provide information on the position on the strip which was hit by the radiation. The best spatial resolution obtained so far is better than $25\mu\text{m}$ and may be improved to less than $10\mu\text{m}$ by reducing the noise of the electronics. The main disadvantage of superconducting strip detectors is their long response time of $1\mu\text{s}$ (for a 1mm long strip) and the much longer dead-time of 1ms . Due to the width of our strips which was larger than $3\mu\text{m}$ we were not able to detect 5.9keV -X-rays. Therefore, narrower strips of widths $w < 1\mu\text{m}$ are required.

We developed a model based on the one-dimensional heat-conduction equation, with which we are able to simulate very well the propagation of a normal-conducting zone in a current-biased superconducting strip. From fits of our model to the experimental data, the thermal coupling between film and substrate and the renormalized value of the Lorenz number can be deduced.

Concerning strip fabrication we have shown that superconducting Ta films on sapphire substrates kept at room-temperature may be obtained by means of the laser-ablation process. The quality of the films depends strongly on the energy density of the laser radiation on the target. A minimum threshold of $200\text{ J/pulse}\times\text{cm}^2$ is required in order to favour the nucleation of Ta in the attempted α -phase. The surface of the films is plagued by sub-micron diameter particles.

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