THERMAL CONDUCTIVITY OF NIOBIUM DISELENIDE

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x

LIST OF SYMBOLS

- $\chi_{\ell m}^k$ - $\chi = K$ - Thermal conductivity
 = W - Thermal resistivity
- k = N - Normal state
 = S - Superconducting state
- ℓ = e - Electron carriers
 = g - Phonon carriers
 = T - Total
- m = e - Electron scatterers
 = i - Impurity scatterers when $\ell = e$
 Non-electron scatterers when $\ell = g$
 = g - Phonon scatterers
 = B - Boundary scattering
 = d - Dislocation scattering
 = p - Point defect scattering
 = u - Umklapp scattering
- T - Superconducting transition temperature
- ρ_0 - Residual electrical resistivity
- ρ - Temperature dependent electrical resistivity
- ℓ - Mean free path
- ℓ - Distance between thermal contacts
- Δ - BCS energy gap parameter
- ΔT - Temperature difference between thermometers

Thermal conductivity of niobium diselenide

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Thermal conductivity measurements have been used to study the temperature dependence of the various scattering mechanisms present in single crystals of NbSe₂ in a temperature range from 1 to 10K. Phonons are found to contribute a substantial fraction of the conduction in this temperature range and, in general, the results agree well with standard theoretical models for three-dimensional metals. The presence of van der Waals boundaries between the layers introduces no significant boundary scattering for phonons. In addition to the usual electron and point defect scattering terms which are present in the phonon conductivity, a term which has a linear temperature dependence was found. This linear term may be caused by stacking fault scattering centers. A linear field dependence of the electrical magnetoresistance is reported and it is suggested this term may arise from magnetic breakdown at energy gaps in the Fermi surface induced by charge-density waves as suggested by Overhauser. Finally it is noted that the ratio of the superconducting to normal state thermal conductivity is consistent with a BCS-like superconducting energy gap.

INTRODUCTION

Superconductivity

When a solid undergoes the normal to superconducting transition, some of its physical properties undergo a dramatic change. This transition is primarily electronic in nature and arises from pair correlations of the electrons (1). The properties of the lattice are essentially unchanged from the normal state properties.

Since complete discussions of the theory of superconductivity by Bardeen, Cooper, and Schrieffer (BCS) and its extensions have been given elsewhere (1,2,3), only the relevant ideas will be discussed here. In the BCS model, the superconducting ground state of a solid consists of a coherent mixture of normal state wave functions in which the states are occupied in pairs rather than as independent particles. The mixing of these Bloch states arises from the coherent exchange of virtual phonons by the electrons and the lowest energy state for the system occurs when electrons with wave vector $\vec{k}$ and spin up are exactly paired with electrons of opposite wave vector and spin down.

This highly correlated superconducting phase of the electron system has been compared to a gigantic molecule (2) in that the electron gas behaves as a unit rather than as individual particles. This highly ordered state leads to a

system of lower entropy than the normal state. In fact, using the two-fluid model (discussed later), the electrons in the superconducting ground state are perfectly ordered and thus have zero-point entropy. This means any electrons participating in the superconducting phase do not participate in any random scattering events because this would randomize the system and increase the entropy to a finite value. This accounts for the fact that superconducting electrons experience no dc electrical resistance and are unable to carry a heat current even though there is a transport of electrons.

Another consequence of this cooperative ordering phenomena is a lowering of the energy of the system. There is a minimum amount of energy required to break one of these superconducting pairs and create an excitation of higher energy. This minimum amount of energy, $E_g = 2\Delta(T)$, accounts for the experimentally observed gap in the superconducting excitation spectrum. The parameter $\Delta(T)$ is thus the energy per particle needed to create an excitation, and is related to the strength of the electron-phonon interaction in the system.

A very important property of a superconductor is that a sufficiently strong magnetic field will destroy superconductivity. It is this property that allows one to take normal state data on a superconductor below the superconducting transition temperature, T_c . In an applied magnetic field a

superconductor exhibits perfect diamagnetism which arises due to circulating surface currents which flow to cancel the flux density inside. For Type-I superconductors, such as most superconducting elements, as the strength of the applied magnetic field is increased, the shielding currents also increase until at some point the critical current of the superconductor is reached and it reverts to its normal state. On the other hand for Type-II superconductors, such as niobium, vanadium, and most superconducting alloys and compounds, the sample goes into the mixed state as the applied field is increased at a lower critical field H_{c1} . The mixed state, which consists of a superconducting matrix with quantized vortices threading through the material, has a lower free energy than it would have if it were completely diamagnetic. As the applied field is increased above H_{c1} , however, the normal regions, called fluxoids, become more closely spaced until finally at an upper critical field, H_{c2} , they merge and the material becomes normal.

Thermal Conductivity

The thermal conductivity of solids at low temperatures is complicated because there are at least two carriers of heat, electrons and phonons, and there are a variety of scattering mechanisms for each of the carriers. For high purity metals most of the heat is carried by electrons, whereas in

impure metals or alloys, a substantial proportion of the thermal conductivity is due to the phonons. For example, the phonon conduction in copper at low temperatures is negligibly small compared to the electronic conduction until about 1% zinc is added (4). At 5% of zinc, about one half of the total conduction is due to phonons. For electron carriers the thermal conductivity is limited by two scattering mechanisms, impurities and phonons, which give rise to thermal resistivities W_{ei} and W_{eg} respectively. Phonon conductivity, on the other hand, is limited by five scattering mechanisms. They are, in order of increasing effectiveness as temperature is increased: 1) scattering of phonons at the specimen boundaries, W_{gB} ; 2) scattering by electrons, W_{ge} ; 3) scattering by dislocations, W_{gd} ; 4) scattering by point defects, W_{gp} ; 5) and finally scattering of phonons by other phonons, W_{gu} , which is negligible at low temperatures.

It is usually assumed that the thermal resistivities are additive so that one can treat each scattering mechanism contribution to the thermal resistivity separately and then add all of them to obtain the total resistivity. This approximation is not strictly valid when the different scattering mechanisms have different frequency dependences. This restriction is much more serious for phonon carriers than electronic carriers due to the very temperature dependent phonon spectrum. There is experimental evidence (5,6,7,8) that this

approximation is indeed invalid in some cases but it will nevertheless be used here in the data analysis. Thus the total electronic thermal resistivity, W_e , is given by

$$W_e = W_{ei} + W_{eg}, \quad (1)$$

and the total phonon resistivity at low temperatures, W_g , is given by

$$W_g = W_{gB} + W_{ge} + W_{gd} + W_{gp}. \quad (2)$$

The total heat conductivity is then found by taking the sum of the electronic and phonon conductivities, i.e.,

$$K_T = K_e + K_g, \quad (3)$$

where $K_e = (W_e)^{-1}$ and $K_g = (W_g)^{-1}$.

There have been numerous theories to explain the temperature dependence of the various scattering mechanisms (9,10,11). Generally, these theories are based on some approximation or simplification of the Boltzmann equation. A rigorous calculation of the conductivity, however, requires the solution of a Boltzmann equation involving complicated integrations which have only been solved for the case of quasi-free electrons. If one considers only the temperature dependence of the thermal conductivity, however, it is useful and instructive to recall the kinetic theory formula for the thermal conductivity

$$K = C v \ell / 3, \quad (4)$$

where C is the specific heat of the carriers, v is their velocity and ℓ is their mean free path. This expression, although admittedly crude, has surprising predictive power in the description of the temperature dependence of real metals. Using this analysis, the temperature dependence will be considered for only those scattering mechanisms which are predominant in this experiment.

For electron carriers, in the case of impurity scattering, the mean free path of the electrons should not change with temperature. The velocity will be considered a constant for both electrons and phonons. W_{ei} will then be proportional to $1/C$, so that W_{ei} will be of the form B/T . The scattering of electrons by impurities also gives rise to a constant residual electrical resistivity, ρ_0 , at low temperatures. ρ_0 and W_{ei} are related through the Wiedemann-Franz law (12)

$$\rho_0 / (W_{ei} T) = L, \quad (5)$$

where L is the Lorenz number and is equal to 2.45×10^{-8} ohm-watt/ K^2 .

Unfortunately there is no simple way to explain the scattering of electrons by phonons and the more detailed theories mentioned earlier are needed even to obtain the temperature dependence. The general results of these theories

(9, 10, 11) are that W_{eg} is of the form AT^2 for low temperatures ($T < \theta/10$). For this simple result to apply we must be in the temperature regime where C varies as T^3 . This does not occur for some metals until T is as low as $\theta/50$ or $\theta/100$ (13) and thus W_{eg} will deviate from the T^2 behaviour until T lies in this region.

It should be mentioned here that the Wiedemann-Franz law is not strictly valid if the scattering is inelastic. The quantity $\rho/W_e T$ is less than the Lorenz number at low temperatures and approaches its theoretical value once again only at higher temperatures. Qualitatively, the reason for this deviation is that electrons interact with low-energy phonons at low temperatures and are scattered through small angles. Since the electronic charge is not altered in the interaction, the current flow is not changed very much from the original direction of flow. On the other hand, the interaction with a phonon, even though it results in a small angle of scattering, changes the electron's energy by kT . The temperature of the electron gas is thus reduced and less heat is transported. Hence W_{eg} is increased more effectively than ρ . At high temperatures where large angle scattering is important the Wiedemann-Franz law is valid again.

The mean free path, λ , of the phonons when being scattered by electrons is inversely proportional to the number of electrons with which they can interact. This number is a

fraction kT/E_F of the total number present since only those electrons within kT of the Fermi surface can interact. Thus the number of electrons available for scattering is proportional to T and ℓ varies as T^{-1} . Again it will be assumed that v is constant for phonons. The specific heat goes as T^3 so W_{ge} will be proportional to T^{-2} .

The mean free path of phonons scattering from point defects is proportional to the fourth power of the phonon wavelength (Rayleigh scattering). Now if the dominant lattice wavelength is inversely proportional to the temperature, ℓ is proportional to T^{-4} . In the region where C varies as T^3 we would expect W_{gp} to be of the form PT . However, we cannot really use the dominant wavelength argument when the scattering process has such a rapidly varying frequency dependence and we must integrate the scattering over the whole phonon spectrum. Generally, this does not yield a linear dependence for W_{gp} .

The last important mechanism for scattering phonons in this experiment is that due to dislocations in the crystal. Following a simplified treatment given by Ziman (14) for scattering by the strain field around the dislocation, we find the scattering cross section is proportional to $1/\lambda$ and thus to T . The mean free path therefore is proportional to T^{-1} and W_{gd} will be equal to DT^{-2} .

In these samples then, we would expect the total phonon resistivity to be of the form

$$W = ET^{-2} + DT^{-2} + (\text{point defect term}). \quad (6)$$

Thermal Conductivity of Superconductors

The thermal conductivity of a superconductor can be qualitatively explained using the two-fluid model (15) in which the electrons can be considered to be two interpenetrating fluids - one the superconducting ground state and the other normal state excitations. The superconductive component as discussed before, has zero-point entropy at all temperatures. At T_c the fractional concentration of normal fluid is 1 and decreases monotonically to zero at $T=0$ K. Because the superconducting component has zero-point entropy, it cannot contribute to the heat transport and therefore K_T^S represents a progressively smaller fraction of K_T^N as T approaches 0 K.

By assuming that excitations from the superconducting state behave independently and using the results of the BCS theory of superconductivity, Bardeen, Rickayzen, and Tewordt (BRT) (16) developed quantitative expressions for the thermal conductivity in the superconductive state for different carriers and scattering mechanisms. For electron carriers being scattered by impurities the BRT expression for the

ratio of the thermal conductivity in the superconducting state to that in the normal state is given by

$$K_e^S/K_e^N = 2F_1(-y) + 2y \ln(1+e^{-y}) + y^2/(1+e^y), \quad (7)$$

where $y = \Delta/kT$ and

$$F_1(-y) = \int_0^\infty Z dZ / (1+e^{Z+y}). \quad (8)$$

A plot of K_e^S/K_e^N vs. T is shown in Fig. 1. In many cases phonon scattering of electrons is important in the temperature range below T_c and K_e^S/K_e^N then takes a form similar to that shown in Fig. 2 (17).

Since phonons as well as electrons contribute to the heat conductivity in a superconductor, the ratio of the total conductivities, K_T^S/K_T^N , is not equal to K_e^S/K_e^N . Indeed the phonon conductivity can be quite different in the superconducting state from that in the normal state. BRT have calculated the change in the conductivity at the normal to superconducting state transition for the case of phonons being scattered by electrons Fig. 3. This behavior of the phonon conductivity can be explained qualitatively by using the fact that the superconducting component of the electron gas does not scatter phonons. As the temperature is reduced below T_c more electrons drop into the superconducting ground state providing less scattering centers for phonons thus explaining the rapid increase in phonon conductivity.

These theoretical considerations along with a wealth of experimental evidence leads us to the conclusion that there are basically three types of temperature dependences for the thermal conductivity in the superconducting state. The first type, shown in Fig. 4a, occurs for pure or slightly doped superconductors. In this type the electrons are the major carrier over most of the temperature range below T_c and impurities and phonons are the dominant scattering mechanisms. The second type, shown in Fig. 4b, occurs for superconductors exhibiting a high phonon conductivity in the normal state. In this situation the phonon conductivity dominates, increasing below T_c due to the decreasing number of scattering centers as discussed before. Finally, some superconductors will exhibit a combination of the above effects as shown in Fig. 4c. Although not always the case, the thermal conductivity usually behaves like the first type near T_c and more like the second type at lower values of T/T_c .

NbSe(2)

Niobium diselenide and other Group V dichalcogenides form a very interesting class of layered, Type-II superconductors having relatively high transition temperatures in the vicinity of 7.2 K. Frindt and Huntley (18), using high purity niobium to prepare the crystals, have reported values of T_c as high as 7.4K for NbSe(2).

NbSe(2) crystallizes in hexagonally packed planes consisting of the Nb metal atoms sandwiched between two chalcogenide (Se) planes with trigonal prismatic symmetry. Weak van der Waals type bonds exist between the layers but bonds within each layer are strong, resulting in a highly anisotropic compound. This anisotropy is clearly evident in electrical resistivity measurements both parallel and perpendicular to the C-axis of the crystal (the C-axis is perpendicular to the planes). Edwards and Frindt (19) have taken such measurements as a function of temperature and have found that the anisotropy ratio, $\rho_{||C}/\rho_{\perp C}$, decreases as the temperature decreases. Coleman and Bhandari (20) have measured the electrical conductivity of NbSe(2) single crystals at 4.2K in magnetic fields up to 150 kOe and have measured the upper critical field H_{C2} as a function of angle. They find a value for the ratio of the effective masses for the electron parallel and perpendicular to the layers, $m_{||C}/m_{\perp C}$, of approximately 10. An additional clue that there exists a two dimensional character arises from heat capacity data (21,22). A T^2 term is found in the results suggesting that the bonding is two dimensional in nature. Crystal imperfections or impurities could have a large effect on the two-dimensional nature of these materials.

There have been several experimental investigations that suggest that superconductivity in NbSe(2) (as well as other

layered superconductors) is mainly two dimensional in nature, arising primarily from the electron-phonon interactions within the layers and is supplemented by interactions between the layers. For example, when the layers of TaS_2 are separated by intercalation (23), T_c rises from 0.8 K to approximately 3 K. On the other hand intercalation of NbSe_2 (23) leads to a drop in T_c from 7K to approximately 3K. Other evidence of the two-dimensional character arises from experiments by Frindt (24). T_c measurements on 2H-NbSe_2 as a function of crystal thickness down to approximately 12 angstroms show that T_c drops from 7K in the bulk material to an extrapolated value of about 3K for a single layer. These measurements indicate that interactions between the layers tend to depress superconductivity in the case of TaS_2 but enhance it in NbSe_2 .

One of the most interesting features of NbSe_2 is the charge-density-wave (CDW) transition which occurs at 33 K. Early studies by Lee et al. (25) showed anomalous behaviour in the Hall coefficient, resistivity, and magnetic susceptibility of 2H-NbSe_2 when measured as a function of temperature. At low temperatures, as the temperature drops the Hall coefficient changes from a constant positive value, passes through zero at a temperature T_x , and then becomes negative. The susceptibility rises in a Curie-Weiss manner until temperatures are in the region of T_x (approximately 40K for

NbSe(2) and then it drops slightly before rising again until T_C is reached. The resistivity drops linearly and then falls off as T^n , where n is between 1 and 2, in the region less than T_X until it reaches T_C . This "anomalous" behavior has recently been attributed to the formation of CDW's (26,27). A CDW would introduce a lattice distortion having the same periodicity as the CDW. In general this distortion is incommensurate with the lattice. As the temperature is decreased the distortion will grow in amplitude and frequently the sample will transform to a commensurate period to minimize strain energies. Earlier work (28,29) has shown there is a slight distortion in the lattice parameters of 2H-NbSe(2) below 40K. Moncton (27), using neutron-scattering studies, has shown that TaSe(2) locks in to a commensurate superlattice at 90 K whereas NbSe(2) remains incommensurate to 5 K.

Thermal Conductivity of NbSe(2)

The purpose of this experimental work is to investigate some of the unusual, and certainly anisotropic, behavior of NbSe(2) using thermal and electrical conductivity as our experimental tools. Thermal and electrical conductivity measurements should produce information on the dimensionality of the electron and phonon spectrum by observing the temperature dependences of the carrier scattering mechanisms. The

heat capacity goes as T^3 for three-dimensional material and as T^2 for two-dimensional material. The thermal conductivity, therefore, should be one power of T lower for the two-dimensional case than the three-dimensional case. It is hoped that the experiments will determine whether NbSe(2) shows a completely two-dimensional behaviour indicative of complete isolation between the layers, or does it behave in a three-dimensional manner showing evidence of its anisotropic bonding only in the behaviour of some of the scattering mechanisms. We want to pay special attention to the role of the boundaries between layers as a scattering mechanism. Through detailed analysis of the scattering mechanisms we should be able to determine the dominant phonon modes existing in the temperature range where this compound becomes superconducting. If there exists many phonon modes at low energy they might contribute a large phonon conductivity at low temperatures.

Measurements on another layer-structure superconducting compound, TaS(1.6)Se(0.4), have shown an unusually large fraction of phonon conductivity to the total heat conductivity (30). Furthermore an anomalous $T^{1.6}$ term was found in the temperature dependence of the phonon conductivity of this compound. We would, therefore, also like to determine if this type of behaviour is found in NbSe(2) and if so, can it be attributed to the anisotropic bonding.

Finally, although it is difficult to determine the superconducting energy gap in a sample with so many scattering mechanisms, we hope to get an estimate of the temperature dependence of the energy gap.

EXPERIMENTAL

Sample Preparation

The samples were made, using the method of iodine vapor transport, by Howard Shanks of this laboratory. The Nb and Se were simultaneously reacted and transported in a 5 cm. diameter, 25 cm. long quartz tube. Sample no. 1 had a reaction temperature of 750°C and transport carried out to 715°C. Samples nos. 2 and 3 were reacted at 600°C and transported to 585°C. After 14 days a number of individual single crystals up to 8x8x0.1 mm were obtained at the cool end of the tube. The transported samples were heated to 100°C in a vacuum of 10^{-6} Torr for 48 hours to remove any iodine adhering to them.

The major impurity in samples No. 2 and No. 3 (from the same batch) was tantalum with a concentration of 600 ppm atomic as determined by mass spectrographic analysis. Sample No. 1 contained equal concentrations of tantalum and iron at 100 ppm.

Cryostat

Temperatures from 1 to 10 K were obtained using a cryostat with a standard He(4) refrigerator and an electric control heater. The cryostat was suspended in a metal, nitrogen shielded helium dewar (Cryogenic Associates, Inc.

model SD-13).

The lower section of the cryostat is shown in Fig. 5. It is very important in this experiment to keep heat leaks to the sample and thermometers to a minimum. To minimize the conduction through residual gas the experimental chamber was enclosed in a vacuum can. The pressure in the can, as measured by a Phillips Ionization Gauge at room temperature, was typically 10^{-6} mm of Hg or less. All leads were first thermally anchored at 4.2 K in the He(4) bath, then at the He(4) chamber, again on the sample tail, and finally on the blocks of copper which contain the thermometers or heater. The proper size (area/length (A/l) ratio) must be chosen for the wires going from the sample tail to the thermometers and heater. A small A/l ratio decreases the heat leak but increases the ohmic heating. The compromise chosen was 10 cm of No. 44 manganin wire for all leads except the sample heater current leads which were 20 cm of No. 36 manganin wire. Using the value of the thermal conductivity and electrical resistivity of manganin given by Anderson et al. (31), and assuming a maximum temperature difference of approximately 0.2 K, the total heat conducted to the sample is on the order of 10^{-8} watts. This includes ohmic heating in all leads except the sample heater current leads. The error in the heater current leads goes as the ratio of the lead resistance to the heater resistance which is about 0.5%. A

first-order correction is obtained by connecting one of the voltage leads at the heater and the other on the circuit board where the leads are thermally anchored to the sample tail. This reduces the error to a few parts in 10^4 . Any heat leak to the sample by radiation is reduced by including the heat shield between the sample and the vacuum jacket.

To insure good thermal isolation of the thermometer mounting blocks and the sample heater mounting block they were suspended from the sample tail by Nymo size A nylon string (see Fig. 5). The total heat leak through the strings is less than 10^{-8} watts. The thermal link from the thermometer mounting blocks to the sample consisted of first, a four centimeter length of #12 copper wire silver soldered to the mounting block, followed by approximately two centimeters of #32 copper wire attached to the #12 wire with GE 7031 varnish. The #32 copper wire was attached to the sample by applying a thin coat of GE 7031 varnish to the wire and carefully placing the wire on the sample surface. After the varnish had dried a thin line of varnish was painted completely around the sample at that point to insure good thermal contact to all the layers. The sample itself was mounted by attaching a thin copper tab to one end with GE 7031 varnish, again making sure all the layers were covered, and the copper tab was varnished to the sample tail of the cryostat. This served as the cold sink contact. Likewise, a

copper tab was varnished to the other end of the sample. Thermal contact to the heater mounting block was made by varnishing this tab to a two centimeter piece of #12 copper wire which had been silver soldered to the mounting block (Fig. 6).

Leads for the purpose of electrical resistivity measurements, which were not attached during the thermal conductivity measurements on sample no. 1, were made to the sample by carefully painting on #40 copper wires with Micro-Circuits Co. SC20 silver paint adjacent to the thermal leads. This proved to be an unsatisfactory procedure for several reasons. First of all, the subsequent handling of the specimen made it highly likely that damage would be done to these fragile samples as was the case in sample No. 1. Furthermore, a new measurement of ℓ/A was required thereby increasing the error in the determination of the thermal conductivity values using the electrical resistivity data. Finally, some of the samples were small enough that space would not permit attaching additional leads adjacent to those already present. For these reasons all other samples used the same #32 copper wire for both thermal and electrical contact to the sample. This was done by applying the varnish as before except over only half of the sample. The other half was painted with the silver paint (Fig. 6). This procedure assures a higher probability that the ℓ/A ratio for the

thermal and electrical measurements are the same and therefore reduces the error in the determination of the thermal conductivity. Furthermore, the sample is only handled once reducing the chance of sample damage between runs.

Thermometry

One of the thermometers in this experiment was the same germanium resistor (GR99) used by Finnemore, Stromberg, and Swenson (32). The other, GR65, was recently calibrated by M. S. Anderson from 1 to 77K in 1974. The calibration of GR99 was checked against GR65 before data was taken. Both thermometers were corrected to IPTS68 (International practical temperature scale 1968). The R vs T values were fit to an equation of the form

$$\ln T(\text{calc}) = \sum_N A(N) (\ln R)^{N-1} \quad (9)$$

by minimizing the root mean square deviation (RMSQD), where RMSQD is given by

$$(\text{RMSQD})^2 = \left[\sum_N (T(N) - T(\text{calc}))^2 \right] / (NPTS - 1).$$

It is important that in the final fit $(T - T(\text{calc}))$ should be small for each point and $(dR/dT)_{\text{calc}}$ and $(d^2R/dT^2)_{\text{calc}}$ should be smooth and independent of the order of fit (33). This minimizes any oscillations that would occur from

"overfitting". The data on each thermometer was fit for two overlapping regions and joined smoothly at a common point.

It is well known that germanium resistance thermometers exhibit a large magnetoresistance in high magnetic fields (34). Since it was necessary to use large magnetic fields to drive the sample into the normal state below T_c , the germanium thermometers had to be calibrated in the field values used. Data was taken by controlling at a given temperature with the capacitance thermometer and taking resistance values at 0, 11, 22, 44 and 54.12 kG. The R vs. T values for a given field value were fit exactly as was done in the zero-field case. The magnetoresistance behaviour of GR65 is shown in Fig. 7.

Neuringer et al. (34) found that the magnetoresistance of the germanium thermometers depends upon its orientation in the field. This orientation dependence is small in fields less than 50 kG (less than 10% at 50 kG for the extreme case of a 180° rotation of the thermometer); nevertheless, it is important to keep the thermometers in a fixed position as much as possible. The magnetoresistance was checked before each run in order to determine if a change in the calibration had occurred. Cycling of the apparatus to room temperature and back and complete removal of the cryostat from the magnet showed no noticeable effect in the calibration.

Temperature Control

In order to measure the thermal conductivity of NbSe(2) below the superconducting transition temperature in the normal state it was necessary to apply magnetic fields on the order of 50 kG. Most resistance thermometers have a magnetoresistance thus making it difficult to use them to control temperature in a strong, varying magnetic field. What one would like is a control thermometer that is insensitive to an applied magnetic field. A very good approximation to this idealized case is found in the glass-ceramic capacitance thermometer. A Lake Shore Cryotronics, Inc., CS-400 GR capacitance thermometer was chosen for this work.

Manufacturers specifications give an equivalent magnetic field temperature error of ± 0.001 K in a field to 150 kG. The thermometer is stable to ± 0.002 K per hour.

In this particular experimental arrangement it was noted that there is a difference in capacitance readings between zero and the maximum applied field. This difference corresponds to an equivalent temperature error of approximately 0.006 K at about 54 kG. Rubin and Lawless (35) also noted a similar effect and attributed it to the type of sample holder and its environment and/or unshielded leads. They speculate that both the capacitance and inductance of the leads and/or thermometer surroundings could be affected by the strong magnetic field. At any rate, in this experi-

ment, the magnitude of the error causes a maximum uncertainty in the thermal conductivity of only 0.2% and no attempt was made to correct the phenomenon.

A transient instability problem has been reported (36) in these thermometers initially believed to be due to helium migration into the sensor body. The magnitude of the drift varied for different sensors but was always $\leq 20\text{mK}$ and would finally stabilize in roughly 30 minutes. There has been a wide variety of reports on the performance of these capacitance thermometers ranging from a 100mK drift in a 4 hour period with no apparent stabilization in sight, to no effect at all. This suggests that the instability is highly dependent on the particular sensor used and/or the experimental set up involved. The thermometer used in this experiment showed no signs of drift on the sub-millidegree scale after a period of 30-40 minutes. No conclusive statement can be made for time periods less than this due to the long time constants involved in this experiment.

The capacitance was measured using a simple three terminal ratio transformer capacitance bridge shown schematically in Fig. 8. For an in-depth discussion of this bridge the reader is referred to a paper by Thompson (37), who developed it. Briefly, the center-tapped secondary of the transformer puts out voltages equal in magnitude and 180° out of phase. If the unknown capacitor is smaller than the standard the

full voltage is applied across the unknown and a fraction of the voltage determined by the ratio transformer is applied to the standard. This fraction is determined by changing the ratio transformer until the current in both arms of the bridge are equal and opposite, creating a null reading on the detector. The ratio of the unknown capacitance to the standard is then equal to the ratio of the voltage applied to the standard to the voltage applied to the unknown. This can be read directly from the dials of the ratio transformer.

The transformer for this bridge is wound on a Supermalloy toroid (Arnold Engineering #4T4178-S1-AA) with a 50 turn primary and 100 turn secondary with 10 taps. To insure that the secondary is as completely independent of the primary as possible, the secondary is electrostatically doubly shielded (38). The ratio transformer is a Gertsch, model RT-60 five dial ratio transformer which is accurate to one part in 10^5 . The standard capacitor is a General Radio, type 1409, three terminal, 0.01 microfarad standard capacitor. It is accurate to better than $\pm 0.01\%$ with a temperature coefficient of $+35 \pm 10$ ppm per degree centigrade between 10° and 70°C . To reduce temperature drift the standard was placed in a styrofoam box for thermal insulation. In this arrangement the temperature drift of the standard is approximately $\pm .0625$ K over a period of four days. A locally made 400 Hz oscillator drives the bridge. A dual phase-

sensitive lock-in detector (39), isolated by a Triad G10 (1:35) Geoformer served as a null detector for the bridge. This detector displays both capacitive and quadrature off-balance simultaneously on separate meters. The oscillator and lock-in were tuned to match each other to 2 parts in 10^3 . The least count voltage switched by the Gertsch ratio transformer is 1 microvolt. This corresponds to a capacitance change of 0.1 pf which means the sensitivity of the bridge is $2.5 \mu\text{V}/0.001 \text{ K}$. The bridge will easily detect a change of 0.0005 K which is suitable for control purposes.

The off-balance of the bridge is fed into an attenuator and zero-shifter which in turn feeds a Bristol mv chart recorder. A cam and microswitch arrangement is mounted on the recorder which is used to switch the current incrementally in the manganin heater at a predetermined set point on the recorder. With this temperature controlling arrangement it is possible to control the temperature to within 0.002 K for 2 hours at 4.50 K, and to within 0.0005 K for periods of 10 to 20 minutes. At higher temperatures (7.0 K) it is only possible to control within 0.003 K for two-hour periods.

Measuring Technique

The resistance of the germanium thermometer was measured using the standard four-terminal potentiometric technique. Current was supplied to the thermometers by a Hg cell current

supply which was stable to a few parts in 10^5 over a half hour. The current was always checked before each measurement. Other components of this system consisted of a Guildline type 9180-B potentiometer, an E.S.I. SR1, 1000.01 ohm standard resistor and a Keithley model 148 nanovoltmeter. This system could easily measure 0.01 ohms out of 2000 ohms which corresponds to a change in temperature of 0.00002 K at 4.00 K and 0.0003 K at 10.00 K. Sample heater current was supplied by a Kepco model CK8-5(M) power supply. The sample heater power was measured using the four-terminal dc technique. The components of this system consisted of a Leeds & Northrup K-3 potentiometer and a Keithley 155 null detector.

The magnetic field was supplied by a Westinghouse 50 kG superconducting solenoid with a two inch bore. At 50 kG the field at the center of the solenoid was constant to .1% over a distance of ± 0.6 inches along the axis. The Westinghouse magnet power supply had a current stability of 2.25 ma (which corresponds to 5 parts in 10^5 at 50 amperes) over 8 hours. The magnet current was measured again using the Leeds & Northrup K-3 potentiometer and Keithley 155 nullmeter by measuring the voltage drop across a Leeds & Northrup 0.01 ohm, 100 amp standard resistor. Fifty amp helium vapor cooled current leads (40) were used to carry the magnet current from the magnet to the top of the dewar.

The procedure for calibrating the germanium resistance thermometers in a magnetic field was to control at a given temperature and sweep through the magnetic field values, coming to equilibrium at each field value. When the magnetic field was changed to a new value, heating was introduced due to magnetic field induced eddy currents, and it was necessary for equilibrium to be established each time. The temperature always returned to the same value (within 0.001 K).

Thermal conductivity data was taken simply by measuring the temperature at each end of the sample and the amount of power applied, $\dot{Q}$, where the heat flow is parallel to the layers. The thermal conductivity is then given by

$$K = (\ell/A) \dot{Q} / (T(H) - T(C)) \quad (10)$$

where (ℓ/A) is the ratio of the length between the thermometer contacts to the cross-sectional area of the sample, $T(C)$ is the temperature at the lead nearest the cold sink and $T(H)$ is the temperature at the lead nearest the heater. Thermal conductivity in a magnetic field, which was applied perpendicular to the layers, was taken in the same manner by keeping the field fixed and sweeping through the temperature values. The error introduced in determining the power input to the sample is approximately a few parts in 10^4 . The difference $T(H) - T(C)$ (ΔT) was typically 0.15 K and the error in determining this value was on the order of 2% or less.

The largest source of error in determining the absolute value of K is the measurement of ℓ/A . This rather large error arises from the inability to make sharply defined thermometer contacts to the sample and difficulty in determining the cross-sectional area of the sample due to its small size and irregular shape. Berman (41) has found that the distances between the midpoints of the areas of the leads can be used as ℓ . The lead evidently takes up the average temperature of the section of the sample that it is touching. This, however, assumes that uniform contact is made between the lead and the sample, which may not be the case. Berman (42) has also considered the problem of the failure of equation (10) when the thermometer leads are too close to the cold sink or to the heater. Under these conditions the temperature gradient is not uniform over the sample. His results show that a spacing of a few mm eliminates this sort of problem. These conditions are marginally met in this experiment.

The area of the sample was determined by the following procedure. First the specimen was photographically enlarged along with a transparent scale. The enlarged area on the photograph was then measured using a planimeter. Comparing this measurement with a standard area measured off the scale gave the area of the sample. The sample was then weighed using an analytical balance. From the weight, area and a calculated

density for NbSe (2) of 6.43 gm/cm^3 the thickness could be found. The width and length ℓ were measured under a microscope with a calibrated reticule. The ℓ/A ratios are known probably no better than a factor of two.

Electrical resistivity measurements were taken once again using the four terminal voltage measurement technique. The voltage drop across the sample was measured using the Guildline potentiometer with the Keithley 155 nullmeter. The current was supplied by the Kepco model CK8-5(M) power supply and was measured by determining the voltage drop across a General Radio 100 ohm $\pm 0.025\%$ standard resistor. The potentiometer and null detector system was also used for this measurement. The uncertainty in the voltage measurement was on the order of 0.05 microvolts. Typically the value of the current used was approximately 12-16 milliamps and the voltage drop across the sample on the order of 2 microvolts. The scatter in the data, therefore, is on the order of 2.5% corresponding to roughly 0.1 micro-ohm cm.

RESULTS AND DISCUSSION

The thermal conductivity for samples 1, 2 and 3 are shown in Figs. 9 - 11. The values of the zero field normal state thermal conductivity in the temperature range below T_C are determined by an extrapolation technique using the thermal magnetoresistance data (Fig. 12). This point will be discussed in detail shortly. For sample 1, K_T^S drops substantially below K_T^N for $T < T_C$ as would be expected if electrons were the major carriers of heat. For samples 2 and 3 the ratio of K_T^S/K_T^N is much closer to one near T_C as would be expected if phonons carried a larger fraction of the heat than in sample 1.

The general method of analysis to separate the electron and lattice components of the total conductivity will be to use the electrical magnetoresistance data in conjunction with the Wiedemann-Franz law and the thermal magnetoresistance data to obtain the lattice thermal conductivity. The electronic and lattice components will then, in turn, be analyzed for their different scattering mechanisms.

Electrical Magnetoresistance

The electrical magnetoresistance data for all the samples can be described very well by linear functions in H . Conventionally, electrical magnetoresistance results are

plotted on a Kohler plot of $\Delta\rho/\rho(0)$ vs. $H/\rho(0)$ where $\rho(0)$ is the resistivity in zero field. For most metals this yields a curve nearly independent of the particular sample used or the temperature. Figs. 14 and 15 show the electrical and magnetoresistance respectively on these Kohler plots. As one can see, there are large deviations from a single curve. These variations could arise if temperature, non-stoichiometry or impurity levels vary the carrier concentration appreciably. This seems unlikely though, because samples 2 and 3 come from the same batch. It is more likely that these deviations come from the inability to accurately measure $\rho(0)$. Fig. 14 also shows a comparison of the Kohler plots derived from this experiment and that of Huntley and Frindt, (HF) (43). The functional form of this data is in qualitative agreement but there is a difference in magnitude of approximately one order of magnitude. If the value of HF's $\rho(0)$, however, is increased by roughly a factor of two the curves would be in much better agreement. They have reported in their paper an unusually low value of the room temperature resistivity of $100 \mu\Omega\text{-cm}$. The values of $\rho(300\text{K})$ in this experiment are from $181 - 193 \mu\Omega\text{-cm}$. These values agree well with previous measurements by Van Maaren and Harland (21), $180 \mu\Omega\text{-cm}$; Lee et al. (25), $160 \mu\Omega\text{-cm}$; and Edwards and Frindt (19), $150 \mu\Omega\text{-cm}$. This discrepancy in HF's data with these data and the data of others is probably the cause of the difference in

magnitude on the Kohler plots.

The most significant aspect of the HF Kohler plot is the large deviations shown for those samples which do not exhibit a reversal of the Hall coefficient. HF suggest that this behavior is indicative of some kind of phase transformation occurring which can be correlated with the reversal of the Hall coefficient. Recently this phase transformation has been identified as arising from the formation of charge-density-waves (26,27) in a rather extensive series of electron and neutron scattering measurements. Unfortunately, the CDW does not occur for all samples but seems to occur only for those sufficiently pure to show a reversal in the Hall coefficient. A charge-density-wave transition has been observed (22) in the batch of NbSe(2) from which samples 2 and 3 were taken. Although sample 1 has not been independently measured for the existence of a CDW transformation, we believe that it does have such a transition because the value of $\rho(0)$ is in the same range as samples 2 and 3 and no large deviations are detected on the Kohler plot.

A central point of interest in this work is the nearly linear dependence on H of both the electrical and thermal magnetoresistance of these samples. Theory (44) predicts that for uncompensated metals the transverse magnetoresistance should saturate in high fields if all orbits are closed, and should vary as H^2 for certain directions which show open

orbits. For compensated metals the only predicted high field behavior is again an H^2 dependence. In view of these theories, the linear dependence is a bit surprising. The apparently anomalous linear behavior in the magnetoresistance, however, has also been observed in many other metals, including potassium (45), sodium (46), aluminum (47), indium (48), zinc (49), and even such substances as single crystal graphite (50). Hence it is not an uncommon phenomenon.

A variety of factors might lead to a linear magnetoresistance. For example one can obtain a linear magnetoresistance by selected averaging over various pieces of the Fermi surface or by appropriate anisotropic relaxation. Alternatively, quasiperiodic static fluctuations, Hall voltage fluctuations, magnetic breakdown due to dislocations and finally, Bragg reflection and magnetic breakdown of electron orbits by CDW introduced energy gaps can be used to predict a linear magnetoresistance. Of all the present explanations, the magnetic breakdown hypothesis seems to be the most likely candidate, at least for NbSe(2).

Magnetic breakdown (51) is the transition of an electron across a gap from one Fermi surface orbit to another under the influence of a magnetic field. The condition for magnetic breakdown to occur (52) is

$$\hbar\omega_H E_F / (E_{\text{gap}})^2 > 1 \quad (11)$$

where ω_H , the cyclotron frequency, is the frequency of an electron orbit in a magnetic field. In a field of 10 kOe in NbSe(2), this corresponds to $E_{\text{gap}} < .07$ eV. This magnetic breakdown can give rise to a linear magnetoresistance because the frequency with which an orbiting electron meets the energy gap is just the cyclotron frequency, $\omega_H = eH/mc$, and hence the scattering is proportional to H .

Pippard (53) developed a formulation whereby dislocations and stacking faults cause the necessary modifications of the Fermi surface to take place to allow for magnetic breakdown. He uses a network model of electron orbits which are coupled by magnetic breakdown and extends his discussion of a two-dimensional metal containing dislocations. Quasiparticles that could travel in straight lines on a perfect network are now heavily scattered to other portions of the Fermi surface. Pippard found that the conductivity was extremely sensitive to the dislocations, even though concentrations may be far too low to affect the zero-field conductivity by direct scattering of electrons.

Overhauser (54) has postulated that a charge-density wave can introduce superzone gaps in the Fermi surface and has shown (55) that this is a plausible explanation for the observed linear behaviour in the magnetoresistance of potassium. The CDW induced magnetic breakdown leads to an approximately linear magnetoresistance which would eventually

have to saturate. The size and k-space location of the energy gap is very sensitive to the orientation of the CDW wave vector, $\bar{Q}$, relative to the crystal axes. This would explain the observed dependence of the transverse magnetoresistance on the orientation of $\bar{H}$ relative to the crystal axes, the wide variation of linear slope from sample to sample, and the sensitivity to deformation.

NbSe(2) might have energy gaps arising from either of the above mechanisms. The samples are known to have undergone the CDW transition and it is very likely that they also have a substantial dislocation or stacking fault network. More work on CDW induced Fermi surface distortions and a systematic study of linear magnetoresistance as a function of CDW occurrence and dislocation density in NbSe(2) is necessary before anything definitive can be said about these mechanisms in relation to an explanation of linear magnetoresistance.

Thermal Magnetoresistance

The determination of the normal state thermal conductivity below T_c is complicated by the relatively large thermal magnetoresistance effect (Fig. 16). Fields of over 50 kOe are required to quench superconductivity and the magnetoresistance at these fields is comparable to the total conductivity. To get the normal state thermal conductivity in zero

field some extrapolation technique must be developed.

To determine the appropriate extrapolation procedure data was taken in a temperature range above T_c where the zero field values of the thermal conductivity can be measured. We then assume that if the data in this range can be fit to a function $K(H,T)$, then one could use this function to determine $K(0,T)$ below T_c . In an effort to find a proper fitting function, we first tried the 2-band theory of Sondheimer and Wilson (56) because of the similarity in the shapes of the plots of $\Delta W/W(0)$ of their theory and of this data (Fig. 17). It is apparent from the graph that Sondheimer and Wilson's expression for $\Delta W/W(0)$,

$$\Delta W/W(0) = C(T) H^2 / (1 + D(T) H^2) \quad (12)$$

was not adequate to give a good enough fit to the data for the accuracy required. Therefore, the following empirical form was used:

$$\Delta W/W(0) = A (1/T)^{n_H m} / (1 + B (1/T)^{n_H m}). \quad (13)$$

This expression was fit to the data by varying the parameters n and m to obtain a best least squares fit. Typical values of m and n are 1.33 and 1.60 respectively. The resulting function would produce values of $K(0)$ to within 1% of the original data points. Although this function would be identical to Sondheimer and Wilson's expression at a given tem-

perature if $m=n=2$, no claim is made that it has any physical significance. The expression is used solely because it will fit the data to the above accuracy. Thus using equation (13) and the thermal magnetoresistance data taken below T_C , "extrapolated" values of $K(0)$ may be determined assuming there is no significant change in the H or T dependence from their values in the temperature range above T_C .

Separation of Electron and Phonon Conductivities

The method of separation of the lattice and electronic contributions to the total conductivity used in this experiment was that of Gruneisen and Adenstedt (57) and De Haas and De Nobel (58). This method is based on the assumption that the phonon conductivity, K_g , and the Lorenz number, L , is unaffected by the application of a magnetic field. With these assumptions,

$$K(T,H) = LT/\rho(H) + K_g(T) \quad (14)$$

where $K(H)$ is the field-dependent total conductivity and $\rho(H)$ is the field-dependent electrical resistivity. If the above assumptions are valid then a plot of $K(H)$ against $T/\rho(H)$ should yield a straight line with slope L and intercept $K_g(T)$ (Fig. 18). Thus after obtaining the curve $K_g(T)$, $K_e(T)$ is obtained by subtracting $K_g(T)$ from $K(T)$ (Fig. 19). L typically had a value of approximately 2×10^{-8} ohm-

watt/K² at 7 K and smoothly increased until it assumed its theoretical value at about 4 K.

As mentioned before, this method rests entirely on the validity of the assumption of the constancy of K_g and L in a magnetic field. It was shown by Makinson (59) that the electron-phonon interaction depends solely on the equilibrium constants of the lattice, provided there is no electric current flowing (56). However, more recently Pippard (60) has postulated that there is a gradual weakening of the electron-phonon interaction as the electronic mean free path becomes small compared to the phonon wavelength, i.e. when $q\ell < 1$. This effect has been experimentally observed (61,62) and could possibly effect this method. The experimental data however, show that this affect is not important in these particular samples in the temperature range of interest. The fact that the plot of $K(H)$ vs. $T/\rho(H)$ gives a straight line gives general validity to this method of analysis. The reason the electron-phonon interaction weakening is not important is that for measurements above 2K in these samples, $q\ell$ is much larger than 1 and the effect would not be easily observed. It has been shown theoretically and experimentally (63) that in the high-field limit L is independent of magnetic field. The experimental data in this experiment, again by virtue of a straight line, show that L is indeed independent of H .

Unfortunately sample No. 1 was badly damaged between the time the thermal conductivity measurements and the electrical resistivity measurements were taken, and the electrical resistivity measurements cannot be trusted to represent the same sample as the thermal conductivity measurements. Therefore, the above method of separation cannot be used for this sample. However, using information from the other samples as a guideline, K_g can still be determined using equation (14). The information needed is that $\rho(H) = \rho(0)(1+AH)$, at least for lower field values (Fig. 13). Then equation (14) becomes

$$K(H) = LT/(\rho(0)(1 + AH)) + K_g \quad (15)$$

which can be rearranged to read

$$H = (L/(\rho(0))) (T/(K(H)-K_g)) - 1/A. \quad (16)$$

Now K_g may be determined by finding that value of K_g that will make the plot of H against $T/(K(H)-K_g)$ linear. This plot also yields the values of A and $L/\rho(0)$. Now if it is assumed that at low temperatures (say 3K) that L assumes its theoretical value of 2.45×10^{-8} watt $^{-1}$ cm deg. then $\rho(0)$ may be determined.

Analysis of Electronic Carrier Scattering Mechanisms

Generally, the main scattering mechanisms in a metal for electronic carriers are impurity and phonon scatterers. It would then seem reasonable to assume the electronic thermal resistivity is given by equation (1). If this equation is written in the form,

$$W_e^{NT} = AT^3 + B, \quad (17)$$

one can plot W_e^{NT} vs. T^3 to determine if the data fit a straight line. Fig. 20 shows that the data fit the linear relation very well giving the constants A and B. Impurity scattering clearly dominates the scattering mechanisms in this temperature range.

One can get a rough idea of the mean free path, ℓ , associated with these scattering mechanisms by using equation (4). Sample No. 2 will be used as representative in this estimate. Using the values of the effective mass of $0.78 m_0$ and the Fermi velocity of 0.51×10^8 cm/sec calculated from Graebner and Robbins (64) and the value of the electronic coefficient of specific heat $\gamma = 10 \text{ mJ/mole-K}^2$ from Bevolo and Shanks (22), the mean free path for impurity scattering is approximately $150 \text{ \AA}$. The mean free path for phonon scattering is then given by the expression $2.29 \times 10^{-4}/T^2$ cm. At 4 K this means that the mean free path due to phonon scattering is $1,430 \text{ \AA}$, roughly a factor of 10 greater than the mean free

path due to impurity scattering. Fig. 21 shows a plot of K_{ei} and K_{eg} .

Two different mechanisms can cause the linear term in the electron conductivity, point defect and boundary scattering. Both mechanisms have scattering cross sections which are independent of temperature. There is no easy way to separate the relative magnitudes of these effects, but one might expect that the mean free path for boundary scattering would be on the order of the boundary spacing, i.e. for NbSe(2) on the order of 12 Å. The observed value of 150 Å would then seem much too large. It is generally believed that it is only diffuse scattering from the boundaries that contributes to the resistivity and in this case one would expect that NbSe(2) has a very smooth surface on an atomic scale. It would seem, therefore, that the predominant scattering at the layer boundaries would be specular in nature and have little effect on the resistivity. The work of Price (65) has shown that specular reflection also contributes to the reduction in conductivity if the Fermi surface is non-spherical. This theory not only gives a basis for the existence of a boundary scattering term, but using the theory and making appropriate assumptions of what "bulk" (no layers) conductivity would be, one can actually arrive at a value in fair agreement with the measured conductivity. The argument is, however, far too speculative in nature and will not be pursued further. The

point here is to demonstrate that the boundary scattering of electrons cannot be ruled out as an effective scattering mechanism in these compounds and indeed, might very well have a significant effect.

$$K_T^S/K_T^N$$

It is constructive at this point in the analysis to look at a plot of the ratio of the thermal conductivity in the superconducting state to that in the normal state vs. reduced temperature. Plots of the ratio K_T^S/K_T^N vs. T/T_C for various samples compared with the theoretical ratio from BRT of K_{ei}^S/K_{ei}^N are shown in Figs. 22 and 23. As expected, the data are not in very good agreement with the theoretical prediction. This, of course, is due to the substantial phonon contribution which we have already discussed in the analysis of the last section.

A reasonable procedure then would be to plot $(K_T^S - K_g^N)/K_e^N$. This is also shown in Figs. 22 and 23. Although this produces a curve in closer agreement with the BRT curve, there is still substantial disagreement. The most probable cause for this discrepancy is the increased phonon conductivity in the superconducting state when the phonons are being less effectively scattered by electrons. Figs. 24 and 25 show what this increase, ΔK_g , over the normal state value would have to be to reach agreement with the BRT curve. The

values ΔK_g are much too small to be in agreement with the BRT theoretical calculations of K_g^S/K_g^N . This just means that only a fraction of the total phonon resistivity arises from electron scatterers, the remaining resistivity being due to some mechanisms unaffected by the superconducting transition. Therefore only this fraction will increase according to BRT and thus would account for the small value we have.

To complicate matters further, with such a large phonon contribution to the conductivity, one might expect the electronic conduction to be limited by phonon scattering as well as impurity scattering, especially at higher temperatures. The lower than average values of K_T^S/K_T^N near T_c gives some indication that this is indeed the case. In view of the large number of variables involved no attempt will be made to calculate K_e^S/K_e^N . Rather we will assume that K_e^S/K_e^N obeys BRT and calculate the phonon conductivity from this assumption. If the resulting phonon conductivities have well-behaved temperature dependences, then one might conclude that the assumption was justified.

Analysis of Phonon Carrier Scattering Mechanisms

The separation of the phonon carrier resistivity into the components for separate scattering mechanisms could use the same approach as was used for the electronic carrier resistivity. Because there are a large number of possible

scattering mechanisms, however, this method would require too many arbitrary constants so the approach was abandoned in favor of the following procedure.

From the experimental raw data and from previous analysis four quantities are known at this point: the total conductivity in the normal state, K_T^N ; the normal state conductivity due to electronic carriers, K_e^N ; the normal state conductivity due to phonon carriers, K_g^N ; and the total conductivity in the superconducting state, K_T^S . In addition, from previous analysis the temperature dependence of the term K_e^N , the normal state electronic conductivity due to phonon scattering, K_{eg}^N , and the normal state electronic conductivity due to impurity scattering, K_{ei}^N , are known. In order to use superconducting theory in our analysis it will be necessary to determine an energy gap, Δ . It will be assumed at this point that the energy gap for NbSe(2) in the planar direction is given by the BCS value of $\Delta = 1.76kT$. This assumption may be justified by the far-infrared transmission measurements of Clayman and Frindt (66) who found the gap in the planar direction to be the BCS value to within their experimental uncertainty. They also found good agreement with the BCS prediction of the temperature dependence of the gap energy and surprisingly found that the anisotropy of the gap was less than 6%.

For impurity scattering, the value of the ratio of the electronic conductivity in the superconducting state to that in the normal state, K_{ei}^S/K_{ei}^N , can be obtained from the BRT theory. Also, for phonon scattering, the superconducting to normal state ratio, K_{eg}^S/K_{eg}^N , can be determined by the theory by Tewordt. Armed with these values and the normal state values of the conductivity, the superconducting values of the electronic conduction for both phonon, K_{eg}^S , and impurity, K_{ei}^S , scattering can be determined. Now the total electronic conductivity can be determined using the relationship

$$(K_e^S)^{-1} = (K_{eg}^S)^{-1} + (K_{ei}^S)^{-1} \quad (18)$$

Using this value of K_e^S in the equation

$$K_T^S = K_e^S + K_g^S \quad (19)$$

yields the value of the total phonon conductivity in the superconducting state. The equations for the phonon resistivities in the normal and superconducting states are,

$$(K_g^N)^{-1} = (K_{ge}^N)^{-1} + (K_{gi}^N)^{-1} \quad \text{and} \quad (20)$$

$$(K_g^S)^{-1} = (K_{ge}^S)^{-1} + (K_{gi}^S)^{-1} \quad (21)$$

where K_{ge}^S and K_{ge}^N are the phonon conductivities in the superconducting and normal states respectively limited by electron scattering, and K_{gi}^S and K_{gi}^N represent all other scattering mechanisms. The quantities on the left-hand side of these

equations are known at this point in the analysis. Now since the transition from the normal to the superconducting state basically affects just the electrons in the crystal, it is reasonable to assume that only the K_{ge} terms will be affected in such a transition and that K_{gi} will remain essentially unchanged (67), i.e. $K_{gi}^S = K_{gi}^N$. This assumption along with the BRT relation between K_{ge}^N and K_{ge}^S gives the needed 4 equations for the 4 unknowns. Equations 20 and 21, therefore, may be solved to find K_{ge}^N and K_{gi}^N . This entire procedure, which is outlined in Fig. 26, needs to be repeated for several different temperatures to give the temperature dependences of K_{ge}^N and K_{gi}^N .

A plot of K_{ge}^N vs. T^2 is shown in Figs. 27 - 29 for the three samples. This plot yields a reasonably straight line showing good agreement with the standard three-dimensional temperature dependence of phonons scattering from electrons as given by eq. 6.

For many of the scattering mechanisms it is difficult to evaluate the absolute value of the resistivity because the number of scattering centers is unknown. For the scattering of phonons by electrons, however, a calculation can be made (61) using the inverse attenuation coefficients obtained by Pippard (68) as mean free paths for longitudinal and transverse acoustic waves. Using this formalism, the lattice conductivity of any metal may be represented on a universal

curve of $(K_{ge}^N / T \rho_0) (b / A \theta_D)^{2/3}$ vs. $(T / \rho_0) (b n \theta_D)^{2/3 1/3}^{-1}$, where ρ_0 is the low temperature electrical resistivity of the sample, b is the valence, A is the gram atomic weight, θ_D is the Debeye temperature, and $n = d/A$, d being the mass density of the metal.

This universal curve, shown as A in Fig. 30, applies only to a metal having a spherical Fermi surface. For non-spherical Fermi surfaces, it is possible that the electron-transverse phonon interaction is stronger (62). Curve B in Fig. 30 shows the lattice conductivity if the attenuation of longitudinal and transverse phonons were identical at high temperatures (69). Fig. 30 also shows the data of two samples of NbSe(2) and various other metals (69). With the exception of the NbSe(2) samples all other data are corrected for the deviation from free electron behaviour by multiplying K_{ge}^N by the ratio α_{exp}/α_{th} , the measured acoustic attenuation divided by the value calculated on the basis of the free electron model of Pippard. This could not be done with NbSe(2) samples because there are no acoustical attenuation data. However, a rough idea of the acoustical attenuation can be obtained by using the thermal conductivity data to obtain a mean free path and thus calculate an α_{exp} . α_{exp} obtained in this way is only the value in the direction parallel to the planes and not an average over crystal orientation as is the case for the other metals. The values of α_{exp}/α_{th}

for NbSe(2) range from 0.55 to 0.97 as compared to a range of 0.66 to 1.4 for the other metals. Thus the NbSe(2) plot is probably not affected much (certainly no more than a factor of two) by not having acoustical data at hand. The value of the valence of NbSe(2) was chosen as one, this being consistent with the high temperature Hall coefficient measurements by Huntley and Frindt (43).

The NbSe(2) data of Fig. 30 are in much better agreement with curve B, both in slope and magnitude, than curve A. The data also compares better with other transition metals than the non-transition metals shown. It is reasonable to assume then, that the NbSe(2) data are consistent with strong scattering of transverse phonons by electrons. Using Bevelo and Shank's (22) value of the lattice coefficient of specific heat, $\beta = 0.686 \text{ mJ/mole-K}^4$ and an average phonon velocity $v = 2 \times 10^5 \text{ cm/sec}$ from the phonon dispersion curves of Wakabayashi and Smith (70), and again using sample No. 2 as representative, the mean free path for longitudinal and transverse phonons scattered by electrons becomes $1.357 \times 10^{-3}/T \text{ cm-K}$ (Fig. 31).

The remaining conductivity, K_{gi}^N , needs further analysis to separate it into various possible scattering mechanisms. Fig. 32 shows K_{gi}^N vs. T along with a smoothed curve drawn through the data points. Fig. 33 shows that these smoothed values multiplied by T and plotted against T^2 yield a reason-

ably straight line. This implies that the resistivity can be written in the form

$$\rho_{gi}^N = AT^{-1} + PT. \quad (22)$$

The dominant contribution to the resistivity in equation 22 is the T^{-1} term; the linear term is an order of magnitude smaller. The linear term can possibly be identified with the scattering of phonons by point defects. If this is the case, the scattering is expected to be small in this temperature range (71). However, recalling the approximations made in arriving at the linear temperature dependence, we cannot say for certain that this term is due to point defect scattering, but may be due to some other combination of scattering terms which only give a small contribution to the total scattering in this temperature range.

The T^{-1} term is somewhat anomalous in the sense that this type of temperature dependence is not usually found in typical metals. Klemens (72) has suggested that such a term would be present when considering the scattering of phonons from stacking faults. It is highly likely that in a crystal such as NbSe(2), stacking faults would be numerous and thus probably contribute to the thermal resistance. Using Klemen's theory we can estimate the number of stacking faults/layer that would give the observed thermal resistivity. Assuming a value of 2 for the Grueneisen parameter, we

arrive at a value of $1.16 \times 10^{-4} T^2$ for the reflection probability and a value of about 0.037 stacking faults/layer. This value of the stacking faults/layer is roughly 8 times larger than Kemp et al. (73) found in strained alpha-brass, and would therefore, seem to be a very reasonable result. The effect is, however, rather large, being on the same order as the lattice resistivity due to electron scatterers. This is in direct conflict with the belief (11) that the effect of stacking faults on long wave phonons would be very small.

Although this T^{-1} term most probably arises from some kind of dislocation or defect structure, the theory of phonon scattering from dislocations is not sufficiently developed (especially for anisotropic crystals) to make any definitive statements as to what exactly causes this term.

SUMMARY

NbSe(2) was found to be a relatively poor conductor of heat compared with most other metals. Fig. 34 shows the total thermal conductivity of NbSe(2) compared with some typical metals and alloys (74,75). This poor conduction arises primarily because in NbSe(2) the energy bands are narrow giving rise to small electron carrier velocities and also because the mean free path is relatively low. The phonons, then, contribute a relatively high proportion of the total heat conductivity, on the order of $1/3$ of the total. A comparison of the total phonon conductivity with some pure and impure metals (76,77,78) is given in Fig. 35.

Making no assumptions about the temperature dependence of the phonon conductivity limited by electron scattering we find it to be in accordance with the usual theoretical prediction for three-dimensional materials. At wavelengths of hundreds of Angstroms the lattice modes are not confined to in-the-plane vibrations in this temperature range. Analysis of this electron scattering term shows that NbSe(2) behaves like other transition metals in that there is a strong scattering of transverse phonons by electrons. We find no evidence of a boundary scattering term in the phonon conductivity as might be expected in this temperature range where the dominant phonon wavelengths are much greater than the

interlayer spacing. The phonons are relatively insensitive to the presence of the layer boundaries.

A rather important result of this experiment is the appearance of a rarely seen linear temperature dependent term in the phonon conductivity. Following the idea of Klemens, we suggest that this term may be due to the effect of stacking fault scattering of phonons. These materials are particularly susceptible to faulting and the magnitude of the mean free path is consistent with a reasonable expectation for fault density (73).

The temperature dependence of the electronic thermal conductivity of NbSe(2) also behaves according to normal three-dimensional theoretical analysis. We find that impurities, phonons, and possibly boundaries are the dominant scattering mechanisms.

A central feature of the magnetic field dependence of the transport properties is the discovery that the electrical magnetoresistance varies linearly with the applied magnetic field. Overhauser (54) has pointed out that magnetic breakdown at the superzone gaps caused by charge-density waves can lead to such a linear behaviour. The NbSe(2) samples are known to have a CDW transition so a linear term might be expected. There are other mechanisms that could give the same temperature dependence, however, and a systematic study of the occurrence of a linear magnetoresistance term as a func-

tion of CDW formation needs to be undertaken. The CDW disappears for samples with resistivities greater than $10 \mu\Omega\text{-cm}$ (43) and the linear term should not be present if CDW's are the cause.

Because there were many different scattering mechanisms present in NbSe(2) we were unable to uniquely determine the value of the superconducting energy gap. However the assumption that the gap is similar to Clayman's value gives phonon conductivities with very reasonable interpretations and which are consistent with the other results of the experiment. We therefore believe that the gap is not too different from Clayman's findings.

Finally, it should be mentioned that although the interpretation of the data in this experiment is not necessarily unique, it is self-consistent and agrees well with other independent data taken on NbSe(2) and with theoretical values of certain parameters.

In broad outline, the thermal conductivity is found to be consistent with a normal three-dimensional metal with a high density of stacking faults, with charge-density-waves and with a BCS value of the superconducting energy gap.

FIGURES

Figure 1. The BRT ratio of the electronic thermal conductivity in the superconducting state to that in the normal state when impurity scattering is predominant

Figure 2. Theoretical curve of the electronic thermal conductivity in the superconducting state to that in the normal state when phonon scattering is predominant

Figure 3. The BRT ratio of the lattice thermal conductivity in the superconducting state to that in the normal state when electronic scattering is predominant

Figure 4. Basic types of thermal conductivity in superconductors: a) pure superconductors, b) superconductors with a high impurity concentration, c) the intermediate situation

Figure 5. He(4) cryostat and sample tail

Figure 6. Blow-up of sample mounting area showing connections to the sample

Figure 7. Magnetoresistance of germanium thermometer GR65

Figure 8. 400 Hz capacitance bridge

Figure 9. Thermal conductivity data for sample 1.
Normal state values below T_C are ex-
trapolated using the procedure in the
text

Figure 10. Thermal conductivity data for sample 2.
Data shows an indication of a higher
phonon contribution to the total con-
ductivity than in sample 1

Figure 11. Thermal conductivity data for sample 3

Figure 12. Representative example of the thermal conductivity data taken at various applied magnetic field strengths

Figure 13. An example of the linear electrical magnetoresistance shown by all the samples

Figure 14. A Kohler plot of the electrical magneto-resistance for NbSe(2) samples. The data are compared with the results of Huntley and Frindt. Temperatures were between 7 and 9 K for the different samples

Figure 15. A Kohler-like plot of the thermal magneto-
resistance for NbSe(2)

Figure 16. Thermal magnetoresistance effect in NbSe(2)

Figure 17. Comparison of the experimental results of the thermal magnetoresistance of NbSe(2) with the theory of Sondheimer and Wilson. Sample No. 1 is used as an example. All the other samples show similar results

Figure 18. $K(H)$ versus $T/\rho(H)$ for sample 2 at various temperatures. The error in the determination of K_g from this graph is approximately 2-3%

Figure 19. Representative example of the electronic and lattice components of the normal state thermal conductivity

Figure 20. Plot of $W_e^N T$ versus T^3 to determine the various scattering contributions. A straight line fit determines the coefficients A and B in $W_e^N = AT^2 + B/T$

Figure 21. Plot of K_{ei}^N and K_{eg}^N . Values represent the thermal conductivity when only the implied scattering mechanisms are present

Figure 22. The ratios of the total conductivity in the superconducting state to that in the normal state for sample 1 are shown by circles. The triangles show the results of subtracting the normal state phonon contribution

Figure 23. The ratio of the total conductivity in the superconducting state to that in the normal state for sample 2. Higher values than Fig. 22 indicate a larger phonon contribution

Figure 24. Correction to K_g necessary to bring the data of Fig. 22^g into agreement with the BRT theoretical curve

Figure 25. Correction to K_g necessary to bring the data of Fig. 23 into agreement with the BRT theoretical curve

Figure 26. Outline of the procedure whereby theoretical superconducting results are used to determine values of κ_{ge}^N and κ_{gi}^N .

FOR A GIVEN T:

$$\textcircled{K_T^N} = \textcircled{K_e^N} + \textcircled{K_g^N} \quad (1)$$

$$\textcircled{K_T^S} = K_e^S + K_g^S \quad (2)$$

$$(K_e^S)^{-1} = (K_{eg}^S)^{-1} + (K_{ei}^S)^{-1} \quad (3)$$

$$(K_g^N)^{-1} = (K_{ge}^N)^{-1} + (K_{gi}^N)^{-1} \quad (4)$$

$$(K_g^S)^{-1} = (K_{ge}^S)^{-1} + (K_{gi}^S)^{-1} \quad (5)$$

RELEVANT EQUATIONS

(PREVIOUSLY CALCULATED
AND MEASURED KNOWNS
ARE CIRCLED)

PROCEDURE FOR FINDING K_{ge}^N AND K_{gi}^N :

1. $\textcircled{(K_e^N)^{-1}} = (K_{eg}^N)^{-1} + (K_{ei}^N)^{-1} = \textcircled{A}T^2 + \textcircled{B}T^{-1} \rightarrow K_{eg}^N, K_{ei}^N$
2. TO FIND K_{eg}^S , USE Twordt, $K_{eg}^S / K_{eg}^N = N' \rightarrow K_{eg}^S$
3. TO FIND K_{ei}^S , USE BRT, $K_{ei}^S / K_{ei}^N = N \rightarrow K_{ei}^S$
4. USE VALUES OF K_{eg}^S AND K_{ei}^S IN EQ.(3) $\rightarrow K_e^S$
5. TO FIND K_g^S , USE K_e^S IN EQ.(2) $\rightarrow K_g^S$

<p>We now have 2 eqs., (4) and (5) and 4 unknowns, $K_{ge}^N, K_{ge}^S, K_{gi}^N, K_{gi}^S \rightarrow$ need 2 more equations</p>
--

6. 1ST EQUATION NEEDED IS BRT, $K_{ge}^S / K_{ge}^N = R$
7. 2ND EQUATION NEEDED IS $K_{gi}^N = K_{gi}^S$
8. SOLVE 4 EQUATIONS, 4 UNKNOWN $\rightarrow K_{ge}^N, K_{gi}^N$
9. INCREMENT T AND RETURN TO STEP 1

Figure 27. Plot of K_{ge}^N versus T^2 for sample No. 1

Figure 28. Plot of κ_{ge}^N versus T^2 for sample No. 2

Figure 29. Plot of K_{ge}^N versus T^2 for sample No. 3

Figure 30. The lattice thermal conductivity of metals, normalized to produce a universal curve A valid for all metals in the free electron approximation. Curve B assumes equal mean free paths for longitudinal and transverse phonons for a ratio of longitudinal/transverse velocity of about 2.2

Figure 31. K_g for various scattering mechanisms
when only those scattering mechanisms
are present

Figure 32. The thermal resistivity due to scattering mechanisms other than electrons

Figure 33. Plot of $W_{gi}^N T$ versus T^2 to determine the various scattering contributions. The straight line fit gives values of A and P of equation 22 for sample 2.

Figure 34. Comparison of the total thermal conductivity in the normal state of NbSe(2) with some typical metals and alloys

Figure 35. Comparison of the total phonon conductivity of NbSe₂ with some pure and impure metals

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APPENDIX

The table which follows presents the fit constants for the equation representing the total normal state thermal conductivity for the three NbSe(2) samples,

$$K_T^N = K_e^N + K_g^N =$$

$$1/(AT^2 + B/T) + 1/(aT^{-2} + b/T + cT).$$

Table 1. Normal state fit constants

	A cm/K-mW	B cm-K ² /mW	a cm-K ³ /mW	b cm-K ² /mW	c cm/mW
Sample 1	0.0000558	0.0992	0.3457	0.1397	0.0014
Sample 2	0.0001390	0.1520	0.6289	0.1407	0.0009
Sample 3	0.0003500	0.2877	1.3500	0.3050	0.0020

NbSe₂
Sample #1
H = kOe

<u>T(K)</u>	<u>K (mW/cm-K)</u>	<u>T(K)</u>	<u>K (mW/cm-K)</u>
1.687	8.376	5.371	64.253
1.997	11.822	5.426	64.376
2.131	13.895	5.521	65.653
2.380	17.879	5.780	69.218
2.538	20.418	5.865	70.412
2.745	23.940	5.997	72.314
2.830	25.422	6.159	74.461
2.956	27.302	6.325	76.736
3.082	29.566	6.586	80.556
3.203	31.662	6.797	83.794
3.368	34.409	6.933	85.832
3.549	37.406	7.134	88.903
3.524	36.996	7.290	90.553
3.648	38.967	7.567	92.891
3.904	42.929	7.843	95.605
4.008	44.555	8.046	97.150
4.229	47.998	8.284	98.929
4.347	49.759	8.518	100.484
4.462	51.468	8.754	102.103
4.627	54.044	9.068	104.256
4.758	55.890	9.366	104.944
4.852	56.969	9.587	107.233
4.970	58.857	10.138	109.702
5.052	59.878	10.485	111.611
5.167	61.511	10.682	112.412
5.317	63.638	11.190	115.244
11.543	116.575	11.953	117.929
12.545	119.339		

$$H = 54.12 \text{ kOe}$$

<u>T (K)</u>	<u>K (mW/cm-K)</u>	<u>T (K)</u>	<u>K (mW/cm-K)</u>
2.007	16.057	5.376	45.535
2.329	18.872	5.753	48.749
2.473	19.817	6.064	51.279
2.789	22.400	6.442	54.603
3.060	24.814	6.682	56.861
3.399	27.945	7.147	60.538
3.665	30.261	7.424	62.744
4.010	33.238	7.692	65.047
4.196	34.667	8.244	69.008
4.599	38.250	8.710	73.668
4.823	39.442	9.164	81.771
5.045	42.477		

$$H = 22 \text{ kOe}$$

<u>T (K)</u>	<u>K (mW/cm-K)</u>	<u>T (K)</u>	<u>K (mW/cm-K)</u>
3.576	32.853	4.335	44.422
3.930	38.180	4.526	46.568
4.085	41.186	4.641	48.681
4.832	50.662	6.612	68.362
5.158	54.060	7.021	72.065
5.468	57.268	7.434	75.654
5.696	59.741	7.839	79.781
5.964	62.521	8.318	84.225
6.210	64.980		

H = 33 kOe

<u>T(K)</u>	<u>K (mW/cm-K)</u>	<u>T(K)</u>	<u>K (mW/cm-K)</u>
2.246	20.113	5.946	56.618
2.528	22.910	6.054	57.749
2.839	25.831	6.127	58.693
2.987	27.270	6.204	59.309
3.163	29.019	6.333	60.347
3.283	30.195	6.590	61.666
3.529	32.622	6.632	62.119
3.904	36.173	6.809	62.759
4.298	40.083	7.128	65.068
4.525	42.264	7.682	71.127
4.918	45.841		
5.101	47.487		
5.399	50.129		
5.658	52.905		

NbSe₂

Sample #1

H = 11 kOe

<u>T(K)</u>	<u>K (mW/cm-K)</u>	<u>T(K)</u>	<u>K (mW/cm-K)</u>
5.013	49.590	6.854	78.480
5.253	56.368	7.287	83.545
5.586	64.901	7.878	88.396
5.873	67.876	8.440	91.619
6.131	70.595		
6.445	73.860		

NbSe₂
Sample #2
H = 0 kOe

$$\lambda/A = 37.5 \text{ cm}^{-1}$$

<u>T(K)</u>	<u>K(mW/cm-K)</u>	<u>T(K)</u>	<u>K(mW/cm-K)</u>
1.648	6.216	7.402	65.130
1.774	8.430	7.619	65.810
1.986	10.736	7.789	66.980
2.127	13.215	7.953	68.560
2.310	15.655	7.989	69.263
2.578	19.800	8.192	69.777
2.815	22.770	8.733	72.312
2.967	25.107	9.111	73.960
3.194	27.850	9.559	75.503
3.459	31.366	10.315	76.897
3.767	34.715	10.804	79.788
3.999	37.370	11.690	82.350
4.543	42.530	12.439	83.470
4.829	44.988		
5.040	46.546	<hr/> H = 54.12 kOe <hr/>	
5.189	48.128	2.025	15.416
5.449	49.673	2.210	17.028
5.598	50.433	2.536	18.795
5.626	51.615	2.984	21.542
5.934	53.822	3.292	24.741
6.031	54.647	3.714	27.954
6.193	55.047	4.070	30.665
6.282	55.940	4.393	32.901
6.343	56.297	4.772	35.428
6.528	57.530	5.062	37.574
6.733	59.140	5.375	39.855
6.922	60.250	5.739	42.077
7.130	62.360	6.002	43.855
7.358	62.945	6.365	46.031

<u>T(K)</u>	<u>K(mW/cm-K)</u>
6.652	48.026
7.028	50.381
7.428	52.502
7.815	54.662
8.081	55.942
8.519	58.520

H = 33 kOe

2.508	20.003
3.282	26.578
4.149	33.974
4.731	38.131
5.334	42.075
6.028	48.026
6.545	50.377
7.037	53.907
7.530	57.079

H = 22 kOe

3.878	32.660
4.0435	34.623
4.703	39.924
5.363	44.852
5.813	48.045
6.362	52.089
6.865	55.812
7.465	59.757
8.1029	64.322

H = 10 kOe

<u>T(K)</u>	<u>K(mW/cm-K)</u>
5.576	49.250
5.897	51.497
6.472	55.630
6.794	58.835
7.444	63.070
8.236	67.192

NbSe₂
Sample #3

H = 0 kOe

$\lambda/A = 20.5 \text{ cm}^{-1}$

H = 54.12 kOe

T(K)	K(mW/cm-K)
1.879	3.850
2.091	5.422
2.387	7.587
2.579	8.897
2.798	10.480
3.069	12.317
3.374	14.330
3.643	15.730
3.871	17.190
4.156	18.500
4.474	19.820
4.686	21.000
4.980	22.230
5.205	23.140
5.448	24.000
5.568	24.370
5.831	25.370
5.999	26.077
6.235	26.878
6.420	28.340
6.566	28.660
6.884	29.820
7.310	31.277
7.713	32.750
7.987	34.297
8.161	35.790
8.582	34.740
8.854	35.670

T(K)	K(mW/cm-K)
2.122	8.527
2.377	9.286
2.600	9.99
2.974	11.66
3.356	13.235
3.728	14.773
4.069	15.14
4.545	17.110
4.978	19.025
5.390	20.44
5.622	21.43
6.043	22.47
6.277	23.90
6.646	24.96
7.051	26.27
7.376	27.12
7.723	27.60
8.139	29.013
8.397	29.916

H = 33 kOe

T(K)	K(mW/cm-K)
2.739	11.04
3.365	13.15
4.146	16.64
4.971	19.87
5.585	22.15
6.295	25.37
6.734	26.625
7.219	28.50
7.961	30.49
7.966	29.306

H = 22 k0e

4.273	18.20
4.569	19.46
5.522	23.26
6.020	24.950
6.499	26.67
7.057	28.52
7.788	30.42
8.150	32.60

H = 11 k0e

5.673	24.34
6.327	26.64
6.943	29.18
7.646	31.23
8.152	32.94