Chapter 5 Phonons II Thermal Properties

Plank Distribution

 $< n_{k,p} >$ is the thermal equilibrium occupancy of phonon wavevector K and polarization p



Quantization of Elastic Waves

- -- The energy of lattice vibration is quantized.
- -- The quantum of energy called a phonon, and the quantum number is denoted as *n*.
- -- The elastic waves in crystals are made of phonons.
- -- Thermal vibration in crystals are thermally excited phonons.

The energy of an elastic mode of angular frequency ω is

$$\epsilon = (n + \frac{1}{2})\hbar\omega$$

(27)

when the mode is excited to quantum number *n*, that is, when the mode is occupied by *n* phonons. The term $\frac{1}{2}\hbar\omega$ is the zero point energy of the mode. Consider a situation where we have a set of identical oscillators in thermal Equilibrium,



Compare the energy difference $h\omega$ of the two adjacent quantum states to the thermal excitation energy k_BT

The ratio of the number of oscillators in their *(n+1)th* quantum state of excitation to the number in *nth* quantum state is

 $N_{n+1}/N_n = \exp(-\hbar\omega/ au)$, $au \equiv k_B T$, The Boltzmann factor(3)

 $N_n = N_0 \exp(-nh\omega/\tau)$

The ratio of the number of oscillators in the *Nth* quantum state to the total number of oscillators is

$$\frac{N_n}{\sum\limits_{s=0}^{\infty} N_s} = \frac{\exp(-n\hbar\omega/\tau)}{\sum\limits_{s=0}^{\infty} \exp(-s\hbar\omega/\tau)} .$$
(4)

We compute the average excitation quantum number *n* of an oscillator to be

$$\langle n \rangle = \frac{\sum_{s} s \exp(-s\hbar\omega/\tau)}{\sum_{s} \exp(-s\hbar\omega/\tau)} .$$
(5)

We let $x = \exp(h\omega / k_B T)$, and $\tau = k_B T$, the denominator becomes

$$\sum_{s} x^{s} = \frac{1}{1-x}$$
Then, the numerator is
$$\sum_{s} sx^{s} = x \frac{d}{dx} \sum_{s} x^{s} = \frac{x}{(1-x)^{2}}, \quad (6)$$

$$\langle n \rangle = \frac{x}{1-x} = \frac{1}{\exp(\hbar\omega/\tau) - 1} \cdot Plank Distribution, \quad (7)$$
as the Eq.(2)

Plank Distribution

 $< n_{k,p} >$ is the thermal equilibrium occupancy of phonon wavevector K and polarization p

$$\langle n \rangle = \frac{1}{\exp(\hbar \omega / \tau) - 1}$$

At large T, $\langle n \rangle \sim \text{ linear in T}$ $\langle n \rangle + \frac{1}{2} \sim (k_{\text{B}}\text{T})/h\omega$ $(\langle n \rangle + \frac{1}{2}) h\omega \sim k_{\text{B}}\text{T}$, as in the classical limit

(2)



Fig. 1

The total energy of the phonons in a crystal is by summing over all phonon modes K, p

$$U = \sum_{K} \sum_{p} U_{K,p} = \sum_{K} \sum_{p} \langle n_{K,p} \rangle \hbar \omega_{K,p} , \qquad (1)$$

We take < *n* > as the Plank distribution

indexed by the phonon wave vector **K**, and the polarization **p**

$$\langle n \rangle = \frac{1}{\exp(\hbar \omega / \tau) - 1}$$

$$I = \sum_{K} \sum_{p} \frac{\hbar \omega_{K,p}}{\exp(\hbar \omega_{K,p}/\tau) - 1} .$$

Contributions from phonons only

(2)

Introducing $D_p(\omega)$, the density of phonon modes (states) at polarization p of frequency in the range from ω to $\omega + d\omega$

$$U = \sum_{p} \int d\omega D_{p}(\omega) \frac{\hbar\omega}{\exp(\hbar\omega/\tau) - 1} \quad \text{Let } \mathbf{X} = \hbar\omega/k_{\text{B}}T \quad (9)$$

$$C_{\text{lat}} = k_B \sum_{p} \int d\omega \ D_p(\omega) \frac{x^2 \exp x}{(\exp x - 1)^2} \xrightarrow{\text{where } C_V \equiv (\partial D/\partial T)_V}{(\text{Lattice Heat Capacity}^{(10)})}$$

(1) For the density of states in one dimension : (A) Fixed Boundary Condition

Consider a 1-dimensional line of N + 1 atoms. Considering N = 10. For fixed boundaries, the end atoms s = 0, and 10 are held fixed, i.e. $u_s = 0$.



The vibration takes up the form of *a standing wave* with the atomic displacement

$$u_s = u(0) \exp(-i\omega_{K,p}t) \sin sKa , \qquad (11)$$

The fixed boundary condition requires that at the end, $u_s = 0$ for s = 0 and s = N, For s = 0, Sin sKa = 0.

For s = N, $NKa = n \pi$, $K = n\pi/Na = n\pi/L$, where n = 1 to N-1.



Showing the *N*-1 allowed values of *K*

$$K = \frac{\pi}{L}$$
, $\frac{2\pi}{L}$, $\frac{3\pi}{L}$, ..., $\frac{(N-1)\pi}{L}$. (12)

For
$$K = \pi/L$$
,
 $u_s \propto \sin(s\pi a/L)$ $u_s = 0$ for $s = 0$ and $s = N$ (13)

For $K = K_{max} = N\pi/L = \pi/a$, $u_s = 0$ for each atom , hence this K is not a valid solution

There are <u>N-1</u> allowed values of K. And there are <u>N-1</u> atoms allowed to move.

*** For the one dimensional line, there is one mode for each interval, $\Delta \mathbf{K} = \pi / \mathbf{L}$

*** And the number of modes per unit range of K is L/π for $K < \pi/a$, and it is zero for $K > \pi/a$, thus $D(K) = L/\pi$

(B) Periodic Boundary Conditions :

For an bounded medium, and require the solutions be periodic over a large distance *L*, namely, u(sa) = u(sa + L), where L = NaWe use the traveling wave solution $U_s = u(0) \exp[i(sKa - w_k t)]$

The periodic boundary condition leads to exp (isaK) = exp (isaK + iKL), then exp (iKL) = 1,

$$KL = \pm 2n\pi, \quad K = \pm 2n\pi/L, \text{ n is integer from 0, to N/2}$$
$$K = 0 \quad , \quad \pm \frac{2\pi}{L} \quad , \quad \pm \frac{4\pi}{L} \quad , \quad \pm \frac{6\pi}{L} \quad , \quad \dots \quad , \quad \frac{N\pi}{L} \quad . \tag{14}$$

There is one mode for every interval $\triangle K = 2\pi/L$. For periodic boundary conditions, the number of modes per unit range of *K* is $L/2\pi$, for $-\pi/a < K < \pi/a$, and is zero otherwise.

$$D(\omega) \ d\omega = \frac{L}{\pi} \frac{dK}{d\omega} \ d\omega = \frac{L}{\pi} \cdot \frac{d\omega}{d\omega/dK}$$
(15)
Group velocity v_{α}

Figure 4 Consider N particles constrained to slide on a circular ring. The particles can oscillate if connected by elastic springs. In a normal mode the displacement u_s of atom s will be of the form sin sKa or cos sKa: these are independent modes. By the geometrical periodicity of the ring the boundary condition is that $u_{N+s} = u_s$ for all s, so that NKa must be an integral multiple of 2π . For N = 8 the allowed independent values of K are 0, $2\pi/8a$, $4\pi/8a$, $6\pi/8a$, and $8\pi/8a$. The value K = 0 is meaningless for the sine form, because sin s0a = 0. The value $8\pi/8a$ has a meaning only for the cosine form, because $\sin(s8\pi a/8a) = \sin s\pi = 0$. The three other values of K are allowed for both the sine and cosine modes, giving a total of eight allowed modes for the eight particles. Thus the periodic boundary condition leads to one allowed mode per particle, exactly as for the fixedend boundary condition of Fig. 3. If we had taken the modes in the complex form exp(isKa), the periodic boundary condition would lead to the eight modes with K = 0, $\pm 2\pi/Na$, $\pm 4\pi/Na$, $\pm 6\pi/Na$, and $8\pi/Na$, as in Eq. (14).



Periodic Boundary Condition for a 1-D Lattice



Figure 5 Allowed values of wavevector K for periodic boundary conditions applied to a linear lattice of periodicity N = 8 atoms on a line of length L. The K = 0 solution is the uniform mode. The special points $\pm N\pi/L$ represent only a single solution because $\exp(i\pi s)$ is identical to $\exp(-i\pi s)$; thus there are eight allowed modes, with displacements of the sth atom proportional to 1, $\exp(\pm i\pi s/4)$, $\exp(\pm i\pi s/2)$, $\exp(\pm i3\pi s/4)$, $\exp(i\pi s)$.



Figure 6 Allowed values in Fourier space of the phonon wavevector K for a square lattice of lattice constant *a*, with periodic boundary conditions applied over a square of side L = 10a. The uniform mode is marked with a cross. There is one allowed value of K per area $(2\pi/10a)^2 = (2\pi/L)^2$, so that within the circle of area πK^2 the smoothed number of allowed points is $\pi K^2 (L/2\pi)^2$.

(2) For the density of States in 3D:

Considering N^3 primitive cells each of edge L

 $\exp[i(K_x x + K_y y + K_z z)] \equiv \exp\{i[K_x (x + L) + K_y (y + L) + K_z (z + L)]\}, \quad (16)$

$$K_x, K_y, K_z = 0 ; \pm \frac{2\pi}{L} ; \pm \frac{4\pi}{L} ; \dots ; \frac{N\pi}{L}$$
 (17)

There is one allowed value of K per volume of $(2\pi/L)^3$ or

$$\left(\frac{L}{2\pi}\right)^3 = \frac{V}{8\pi^3}$$
 Allowed value of K per unit volume of K space. (18)

For a volume of a sphere of radius of K, the total number of modes of wavevectors less than K is

$$N = (L/2\pi)^3 (4\pi K^3/3) \tag{19}$$

The density of states for each polarization is

$$D(\omega) = dN/d\omega = (VK^2/2\pi^2)(dK/d\omega) .$$
(20)
= (dN/dK) / (dK/d\omega)

Debye Model for Density of States

In the Debye approximation, the velocity of sound is taken as a constant. v is the velocity of sound, this usually applies for small ω and small K, such as for the acoustic mode for an elastic continuum.

$$\omega = vK , \qquad (21)$$

$$N = (V\omega^3) / 6\pi^2 v^3 \quad \text{after Eq. 19}$$

$$D(\omega) = V\omega^2 / 2\pi^2 v^3 . \qquad (22)$$

In the Debye model, we define a maximu cuttoff frequency ω_D , and a maximum cutoff wavevector K_D

$$\omega_D^3 = 6\pi^2 v^3 N/V \ . \tag{23}$$

$$K_D = \omega_D / \upsilon = (6\pi^2 N / V)^{1/3}$$
 (24)

The thermal energy of phonons in a crystal is

$$U = \int d\omega \ D(\omega) \langle n(\omega) \rangle \hbar \omega = \int_0^{\omega_D} d\omega \left(\frac{V \omega^2}{2\pi^2 v^3} \right) \left(\frac{\hbar \omega}{e^{\hbar \omega/\tau} - 1} \right) , \qquad (25)$$

$\boldsymbol{\omega}$ vs $\boldsymbol{\kappa}$ Dispersion for Monoatomic Lattice



Figure 4 Plot of ω versus K. The region of $K \ll 1/a$ or $\lambda \gg a$ corresponds to the continuum approximation; here ω is directly proportional to K.

Assuming phonon velocity \boldsymbol{v} is independent of polarization for all three direction, We can multiply a factor of three

$$U = \frac{3V\hbar}{2\pi^2 v^3} \int_0^{\omega_D} d\omega \, \frac{\omega^3}{e^{\hbar\omega/\tau} - 1} = \frac{3Vk_B^4 T^4}{2\pi^2 v^3 \hbar^3} \int_0^{x_D} dx \, \frac{x^3}{e^x - 1} \,, \qquad (26)$$

where $\mathbf{x} = h\omega_D / k_B T$

$$x_D \equiv \hbar \omega_D / k_B T \equiv \theta / T \quad . \tag{27}$$

We Define Debye temperature heta

 $\theta = h\omega_D / k_B$

$$\theta = \frac{\hbar v}{k_B} \cdot \left(\frac{6\pi^2 N}{V}\right)^{1/3} , \qquad (28)$$

The total phonon energy is

$$U = 9Nk_BT\left(\frac{T}{\theta}\right)^3 \int_0^{x_D} dx \, \frac{x^3}{e^x - 1} \quad , \tag{29}$$

 $C_V \equiv (\partial U/\partial T)_V$ By differentiating the eq. (26), we get

$$C_{V} = \frac{3V\hbar^{2}}{2\pi^{2}v^{3}k_{B}T^{2}} \int_{0}^{\omega_{D}} d\omega \frac{\omega^{4} e^{\hbar\omega/\tau}}{(e^{\hbar\omega/\tau} - 1)^{2}} = 9Nk_{B} \left(\frac{T}{\theta}\right)^{3} \int_{0}^{x_{D}} dx \frac{x^{4} e^{x}}{(e^{x} - 1)^{2}} .$$
(30)

For large I, x is small, we can expand the integrand, and $Jax C_v = 3Nk_B$, since $U = 3Nk_BT$ in the large T limit.

For low T, x_D is approaching to infinity, thus in Eq. (29),

$$\int_{0}^{\infty} dx \, \frac{x^{3}}{e^{x} - 1} = \int_{0}^{\infty} dx \, x^{3} \sum_{s=1}^{\infty} \exp(-sx) = 6 \sum_{1}^{\infty} \frac{1}{s^{4}} = \frac{\pi^{4}}{15} , \qquad (31)$$
$$C_{V} \approx \frac{12\pi^{4}}{5} \, Nk_{B} \left(\frac{T}{\theta}\right)^{3} \approx 234 \, Nk_{B} \left(\frac{T}{\theta}\right)^{3} , \qquad (32)$$

The Debye T^3 Law for low temp





Fig. 8



Fig. 9 The low temperature heat capacity of solid Argon linear dependence of T^3



Of the allowed volume in **K** space, the fraction occupied by the excited modes is of the order of $(\omega_T / \omega_D)^3$ or $(K_T / K_D)^3$, where K_T is the thermal wavevector defined as $h\omega_T = k_B T$, $\omega_T = v K_T$.

 $h_V K_T = k_B T$, and $h_V K_D = k_B \theta$

Thus the fraction occupied is $(T/\theta)^3$ of the total volume in the *K* space For **3N** modes of energy $K_B T$, we get $U = 3NK_B T (T/\theta)^3$, and $C_V \sim \partial U/\partial T = 12 NK_B (T/\theta)^3$ as the T³ dependence

Einstein Model of Density of States

 $D(w) = N \,\delta(w - w_o) \qquad \text{We have } N \text{ oscillators of identical fequency } w_o,$ $U = N \langle n \rangle \hbar \omega = \frac{N \hbar \omega}{e^{\hbar \omega / \tau} - 1} , \qquad (33)$ $C_V = \left(\frac{\partial U}{\partial T}\right)_V = N k_B \left(\frac{\hbar \omega}{\tau}\right)^2 \frac{e^{\hbar \omega / \tau}}{(e^{\hbar \omega / \tau} - 1)^2} , \qquad (34)$

At high T, **x** is << 1, $C_V = 3N k_B$, as the Dulong and Petit value.

At low T, **x** is >>1, $C_V = exp(-h\omega/k_BT)$



This model usually applies to **optical phonon modes** where ω is nearly independent of *K*.



 $K = K_{\text{max}} = \pi/a$. The lattice constant is a.

This model usually applies to optical phonon modes, where ω is nearly independent of *K*.

Optical and Acoustic Branches of the Dispersion for a Diatomic Linear Lattice





Fig. 11 Experimental C_v vs T/θ for diamond compared to the Einstein Model In dashed line

General Expressions for $D(\omega)$: As the density of states per unit frequency range at a given $\omega(K)$

For phonon frequency between ω and $\omega + d\omega$

$$D(\omega) \ d\omega = \left(\frac{L}{2\pi}\right)^3 \int_{\text{shell}} d^3 K \quad , \tag{35}$$

The integral is extended over the volume of shell in *K* space bounded between two surfaces of constant frequency ω , and constant frequency $\omega + d\omega$

The volume between the constant frequency surfaces ω and $\omega + d\omega$ is a right cylinder of base dS_{ω} , and altitude dK_{\perp}

$$\int_{\text{shell}} d^3 K = \int dS_{\omega} dK_{\perp} \quad .$$

Where dK_{\perp} is the perpendicular distance between two constant frequency surfaces, and dS_w is an element of area of the constant frequency surface of ω in K space.





Fig. 13

The quantity dK_{\perp} is the perpendicular distance between two constant frequency surfaces ω , and $\omega + d\omega$

The gradient of ω , which is $\nabla_{\mathbf{K}}\omega$, is also normal to the surface ω constant, and the quantity

$$\nabla_{\mathbf{K}}\omega|\ dK_{\perp}=d\omega \ ,$$

is the difference in frequency between the two surfaces connected by dK_{\perp} . Thus the element of the volume is

$$dS_{\omega} \ dK_{\perp} = dS_{\omega} \ \frac{d\omega}{|\nabla_{\mathbf{K}}\omega|} = dS_{\omega} \ \frac{d\omega}{v_{g}}$$

where $v_g = |\nabla_{\mathbf{K}} \omega|$ is the magnitude of the group velocity of a phonon. For (35) we have

$$D(\omega) \ d\omega = \left(\frac{L}{2\pi}\right)^3 \int \frac{dS_{\omega}}{v_g} \ d\omega$$
.





The density of states **D** (ω) vs ω for (a) the Debye Solid, (b) the actual solid. The spectrum for the crystal structure starts as ω^2 , but discontinuity develops at the singular points ($v_g = 0$)

Known as the Van Hove Singularity



Harmonic interaction: Two phonon waves do not interact, no thermal expansion. Anharmonic Crystal Interactions

The energy includes higher order terms than the quadratic (harmonic) term.

 $U_3 = A e_{xx} e_{yy} e_{zz}$

Three phonon processes are caused by the third order terms in the lattice potential energy.



By using the Boltzmann distribution, we calculate <u>average displacement</u> < x > as the ratio of the following two terms:

$$\langle x \rangle = \frac{\int_{-\infty}^{\infty} dx \ x \ \exp[-\beta U(x)]}{\int_{-\infty}^{\infty} dx \ \exp[-\beta U(x)]} ,$$

 $\int dx \, x \, \exp(-\beta U) = \beta = 1/k_B T$

$$\cong \int dx \, [\exp(-\beta c x^2)](x + \beta g x^4 + \beta f x^5) = (3\pi^{1/2}/4)(g/c^{5/2})\beta^{-3/2} ;$$

$$\int dx \, \exp(-\beta U) \cong \int dx \, \exp(-\beta c x^2) = (\pi/\beta c)^{1/2} \quad , \tag{39}$$

$$\langle x \rangle = \frac{3g}{4c^2} k_B T \qquad (k_B T)^2 \qquad (40)$$

<x> is linear at higher T, but
<x> is ~0 at very low T



Fig. 15 Lattice constant of solid Argon as a function of temperature

THERMAL CONDUCTIVITY

$$j_U = -K \, \frac{dT}{dx} \; ,$$

 j_U as the flux of thermal energy transmitted per unit time per unit area (41) dT/dx is the temperature gradient

The transfer of thermal energy process is not a straight path, but a *random* diffusion process with frequent collisions.

 $j_U = \langle N_f \rangle c \Delta T$, $\langle N_f \rangle$ is the particle flux = $n \langle v_x \rangle$ in the x direction Here we introduce

 $\Delta T = dT/dx \,\ell_x = dT/dx \,v_x \tau \qquad \ell_x \text{ is the mean free path between collisions}$ $\tau \text{ is the average time between collisions}$

$$j_{U} = -n \langle v_{x}^{2} \rangle c\tau \frac{dT}{dx} = -\frac{1}{3}n \langle v^{2} \rangle c\tau \frac{dT}{dx} . \qquad (43)$$

$$j_{U} = -\frac{1}{3}Cv\ell \frac{dT}{dx} , \qquad \text{let } \ell = v\tau, \ C = nc \qquad (44)$$

$$K = \frac{1}{3}Cv\ell , \qquad Debye's Expression for K (42)$$

Table 2 Phonon mean free paths

[Calculated from (44), taking $v = 5 \times 10^5$ cm/sec as a representative sound velocity. The ℓ 's obtained in this way refer to umklapp processes.]

Crystal	<i>T</i> , °C	C, in J cm ^{-3} deg ^{-1}	K, in W cm ⁻¹ deg ⁻¹	ℓ, in Å
Quartzª	0	2.00	0.13	40
	-190	0.55	0.50	540
NaCl	0	1.88	0.07	23
	-190	1.00	0.27	100

"Parallel to optic axis.



Open end tube



Fig. 16a

Figure 16a Flow of gas molecules in a state of drifting equilibrium down a long open tube with frictionless walls. Elastic collision processes among the gas molecules do not change the momentum or energy flux of the gas because in each collision the velocity of the center of mass of the colliding particles and their energy remain unchanged. Thus energy is transported from left to right without being driven by a temperature gradient. Therefore the thermal resistivity is zero and the thermal conductivity is infinite.

 J_u is finite, dT/dX = 0, thus the thermal conductivity **K** is infinite !

Close end tube



Fig. 16b

Figure 16b The usual definition of thermal conductivity in a gas refers to a situation where no mass flow is permitted. Here the tube is closed at both ends, preventing the escape or entrance of molecules. With a temperature gradient the colliding pairs with above-average center of mass velocities will tend to be directed to the right, those with below-average velocities will tend to be directed to the left. A slight concentration gradient, high on the right, will be set up to enable the net mass transport to be zero while allowing a net energy transport from the hot to the cold end.

Net energy tansfer = Ju is finite

(A) For a three phonon collision process: the total momentum of the phonon gas is not changed by such a collision.

 $\mathbf{K}_1 + \mathbf{K}_2 = \mathbf{K}_3 \qquad \qquad \mathbf{G} = \mathbf{0}, \qquad \text{the N process} \qquad (45)$

The total phonon momentum is conserved.

$$\mathbf{J} = \sum_{\mathbf{K}} n_{\mathbf{K}} \hbar \mathbf{K} \tag{46}$$

(B) *The <u>Umklapp</u> Process*: the three phonon process that caused thermal Resistivity,

 $K_1 + K_2 = K_3 + G$, G is a reciprocal lattice vector, the U process (47)

The energy of phonons \mathbf{K}_1 , \mathbf{K}_2 suitable for umklapp to occur is of the order of $\frac{1}{2}k_B\theta$, because each of the phonons 1 and 2 must have wavevectors of the order of $\frac{1}{2}G$ in order for the collision (47) to be possible.



The N Process

Figure 16c In a crystal we may arrange to create phonons chiefly at one end, as by illuminating the left end with a lamp. From that end there will be a net flux of phonons toward the right end of the crystal. If only N processes $(K_1 + K_2 = K_3)$ occur, the phonon flux is unchanged in momentum on collision and some phonon flux will persist down the length of the crystal. On arrival of phonons at the right end we can arrange in principle to convert most of their energy to radiation, thereby creating a sink for the phonons. Just as in (a) the thermal resistivity is zero.

K = 0



Figure 16d In *U* processes there is a large net change in phonon momentum in each collision event. An initial net phonon flux will rapidly decay as we move to the right. The ends may act as sources and sinks. Net energy transport under a temperature gradient occurs as in (b).

Net energy transfer

The momentum *G*, and the energy are transmitted to the crystal, causing the temperature change.

 J_u = finite, dT/dx = finite.



Figure 17 (a) Normal $\mathbf{K}_1 + \mathbf{K}_2 = \mathbf{K}_3$ and (b) umklapp $\mathbf{K}_1 + \mathbf{K}_2 = \mathbf{K}_3 + \mathbf{G}$ phonon collision processes in a two-dimensional square lattice. The square in each figure represents the first Brillouin zone in the phonon \mathbf{K} space; this zone contains all the possible independent values of the phonon wavevector. Vectors \mathbf{K} with arrowheads at the center of the zone represent phonons absorbed in the collision process; those with arrowheads away from the center of the zone represent phonons emitted in the collision. We see in (b) that in the umklapp process the direction of the x-component of the phonon flux has been reversed. The reciprocal lattice vector \mathbf{G} as shown is of length $2\pi/a$, where a is the lattice constant of the crystal lattice, and is parallel to the K_x axis. For all processes, N or U, energy must be conserved, so that $\omega_1 + \omega_2 = \omega_3$.

Crystal imperfections:

We must consider scattering by crystal boundaries, the distribution of isotopic masses in natural chemical elements, chemical impurities, lattice imperfections, and amorphous structures.

At high $T, \ell \sim 1/T$ At low $T, \ell \sim D$, limited by the sample width

At low T, $C \sim T^3$

 $K \approx C v D \quad . \tag{48}$

K varies as T^3 at low *T*, and displays a maximum feature due to the temperature dependence of ℓ .



Figure 18 Thermal conductivity of a highly purified crystal of sodium fluoride, after



Figure 19 Isotope effect on thermal conduction in germanium, amounting to a factor of three at the conductivity maximum. The enriched specimen is 96 percent Ge⁷⁴; natural germanium is 20 percent Ge⁷⁰, 27 percent Ge⁷², 8 percent Ge⁷³, 37 percent Ge⁷⁴, and 8 percent Ge⁷⁶. Below 5 K the enriched specimen has $K = 0.060 T^3$, which agrees well with Casimir's theory for thermal resistance caused by boundary scattering. (After T. H. Geballe and G. W. Hull.)

In an otherwise perfect crystal, the distribution of <u>isotopes</u> of the chemical elements often provides an important mechanism for <u>phonon scattering</u>. The random distribution of isotopic mass disturbs the periodicity of the density as

This leads to a significant increase in thermal conductivity