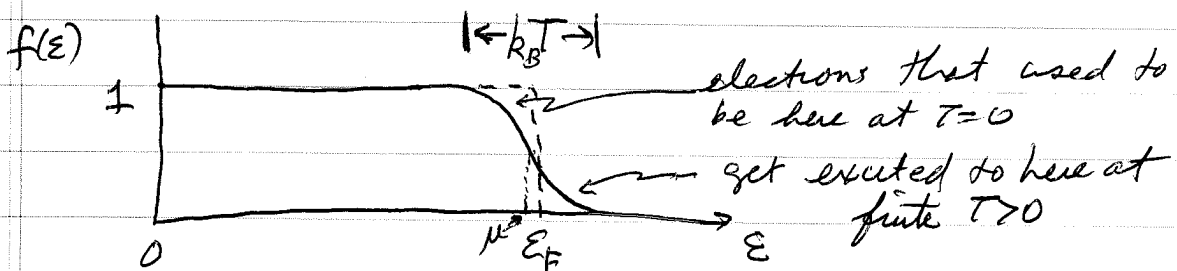


At finite temperature, electrons can absorb thermal energy $k_B T$ from the heat reservoir that keeps the system at temperature T (we are thinking of the canonical ensemble)

But for an electron with $\epsilon \ll \epsilon_F - k_B T$, absorbing $k_B T$ energy would just take it to a state that is already filled. Pauli exclusion principle prevents this from happening.

\Rightarrow The only electrons which can change their state when the system is at finite temperature are those within $\sim k_B T$ of ϵ_F . Only these states can go from a filled to an empty state by absorbing $k_B T$ energy from the heat bath



when $\epsilon = \mu$, $f(\epsilon) = 1/2$

Since $g(\epsilon) \sim \sqrt{\epsilon}$, there are more empty states

within $k_B T$ above E_F than there are within $k_B T$ below E_F .

Since $f(\mu) = 1/2$ is the dividing point ^{in energy} between states which are more likely filled and those that are more likely empty (since $f(\epsilon)$ is also the probability the state is occupied) we see that at finite temperature $\mu(T, n)$ must decrease somewhat from its value E_F at $T=0$.

A more precise calculation (see the "Sommerfeld expansion" in Ashcroft + Mermin pgs 45-47) gives the result (Ashcroft + Mermin Eq (2.78))

$$\mu(T, n) = E_F \left[1 - \frac{1}{3} \left(\frac{\pi k_B T}{2 E_F} \right)^2 \right]$$

↑ gives dependence on T
dependence of μ on n is via E_F

Since $T_F \sim 10^4$ K for most metals, the correction to the chemical potential at finite temperature are

$$|\mu(T) - E_F| \sim \left(\frac{T}{T_F} \right)^2 \sim 10^{-4} \text{ at room temperature}$$

So for most purposes it will be OK to approx $\mu(T) \approx E_F$ even at finite temperature.

Derivation of $\mu(T)$ for $T \ll T_F$

density of electrons is $n = \int_0^{\infty} d\varepsilon g(\varepsilon) f(\varepsilon)$

density of electrons remains fixed as temperature varies

$$\Rightarrow \frac{\partial n}{\partial T} = 0 = \int_0^{\infty} d\varepsilon g(\varepsilon) \frac{\partial f}{\partial T}$$

fermi function $f(\varepsilon) = \frac{1}{e^{x+1}}$ where $x = \frac{\varepsilon - \mu(T)}{k_B T}$

$$\text{so } \frac{\partial f}{\partial T} = \frac{\partial f}{\partial x} \frac{\partial x}{\partial T} = \left(-\frac{\partial f}{\partial x} \right) \left[\frac{x}{T} + \frac{1}{k_B T} \frac{d\mu}{dT} \right]$$

$$d\varepsilon = k_B T dx$$

$$\Rightarrow 0 = \int_{-\frac{\mu}{k_B T}}^{\infty} dx g(\varepsilon) \left(-\frac{\partial f}{\partial x} \right) \left[\frac{x}{T} + \frac{1}{k_B T} \frac{d\mu}{dT} \right]$$

since $\left(-\frac{\partial f}{\partial x} \right)$ is strongly peaked about $\varepsilon = \mu$ and vanishes for $\varepsilon \ll \mu$ or $\varepsilon \gg \mu$,

we can ① extend lower limit of integration to $-\infty$

② approximate $g(\varepsilon) \approx g(\mu) + g'(\mu)(\varepsilon - \mu)$

$$= g(\mu) + g'(\mu) k_B T x$$

then

$$0 = \int_{-\infty}^{\infty} dx \left[g(\mu) + g'(\mu) k_B T x \right] \left[\frac{x}{T} + \frac{1}{k_B T} \frac{d\mu}{dT} \right] \left(-\frac{\partial f}{\partial x} \right)$$

$$\text{Now } \left(-\frac{\partial f}{\partial x} \right) = \frac{e^x}{(e^x + 1)^2}$$

one can check that $\left(-\frac{\partial f}{\partial x} \right)$ is symmetric function of x

$$\text{Therefore } \int_{-\infty}^{\infty} dx \left(-\frac{\partial f}{\partial x} \right) x = 0$$

So we get

$$0 = \frac{g(\mu)}{k_B T} \frac{d\mu}{dT} \int_{-\infty}^{\infty} dx \left(-\frac{\partial f}{\partial x} \right) + g'(\mu) k_B \int_{-\infty}^{\infty} dx \left(-\frac{\partial f}{\partial x} \right) x^2$$

$$\text{where } \int_{-\infty}^{\infty} dx \left(-\frac{\partial f}{\partial x} \right) = f(-\infty) - f(\infty) = 1 - 0 = 1$$

$$\int_{-\infty}^{\infty} dx \left(-\frac{\partial f}{\partial x} \right) x^2 = \int_{-\infty}^{\infty} dx \frac{x^2 e^x}{(e^x + 1)^2} = \frac{\pi^2}{3}$$

$$0 = \frac{g(\mu)}{k_B T} \frac{d\mu}{dT} + \frac{\pi^2}{3} g'(\mu) k_B$$

$$\Rightarrow \frac{d\mu}{dT} = - \frac{\pi^2}{3} \frac{g'(\mu)}{g(\mu)} k_B T$$

$$\text{expect } \frac{g'(\mu)}{g(\mu)} \sim \frac{1}{\mu} \sim \frac{1}{E_F} \text{ at low } T$$

$$\text{so } \frac{d\mu}{dT} \sim \frac{k_B T}{E_F} \sim \left(\frac{T}{T_F} \right) \sim 10^{-2} \text{ is small for all metals}$$

Therefore can approximate

$$\frac{d\mu}{dT} \approx -\frac{\pi^2}{3} \frac{g'(\epsilon_F)}{g(\epsilon_F)} k_B^2 T$$

↑ replaced μ by ϵ_F in argument of g and g' . Corrections to this just give terms higher order in T/T_F

$$\mu(T) = \epsilon_F - \frac{\pi^2}{6} \frac{g'(\epsilon_F)}{g(\epsilon_F)} (k_B T)^2$$

$$\mu(T) = \epsilon_F \left[1 - \frac{\pi^2}{6} \frac{g'(\epsilon_F)}{g(\epsilon_F)} \frac{(k_B T)^2}{\epsilon_F} \right]$$

For free electron gas, $g(\epsilon) = \frac{3}{2} \frac{m}{\epsilon_F} \sqrt{\frac{\epsilon}{\epsilon_F}}$, $g(\epsilon_F) = \frac{3}{2} \frac{m}{\epsilon_F}$

$$g'(\epsilon) = \frac{3}{4} \frac{m}{\epsilon_F} \sqrt{\frac{\epsilon}{\epsilon_F}} \frac{1}{\epsilon}, \quad g'(\epsilon_F) = \frac{3}{4} \frac{m}{\epsilon_F^2}$$

$$\mu(T) = \epsilon_F \left[1 - \frac{\pi^2}{6} \frac{\left(\frac{3}{4} \frac{m}{\epsilon_F^2}\right)}{\left(\frac{3}{2} \frac{m}{\epsilon_F}\right)} \frac{(k_B T)^2}{\epsilon_F} \right]$$

$$\mu(T) = \epsilon_F \left[1 - \frac{\pi^2}{12} \left(\frac{k_B T}{\epsilon_F}\right)^2 \right]$$

$$\text{so } \frac{\mu(T) - \epsilon_F}{\epsilon_F} \sim 0 \left(\frac{k_B T}{\epsilon_F}\right)^2 \sim 10^{-4}$$

⇒ In most situations can ignore T dependence of μ and take $\mu(T) = \epsilon_F$

Specific Heat of Electron Gas

Hand waving argument

Specific heat per volume $c_v = \frac{dU}{dT}$
where $u = \frac{E}{V}$ is energy per volume

When start at $T=0$ and increase temperature to T , it is only the electrons within $k_B T$ of E_F that can absorb $\Delta E = k_B T$ of energy from the thermal bath. So the change in energy per volume is

$$u - u_0 = \Delta u = k_B T \times (\text{number of electrons that absorb energy from heat bath per volume})$$

The ~~most~~ number of electrons that absorb energy is just the number of electrons with $\Delta E = k_B T$ of the Fermi energy. This is $g(E_F) \Delta E = g(E_F) k_B T$
 \uparrow density of electrons per energy per volume

So

$$\Delta u = (k_B T) g(E_F) k_B T = g(E_F) (k_B T)^2$$

$$\text{So } c_v = \frac{d\Delta u}{dT} = 2 g(E_F) k_B^2 T \sim T$$

Specific heat of conduction electrons at $T \ll T_F$ is linear in temperature

and proportional to the density of states at the Fermi energy $g(E_F)$. We will see that many thermodynamic quantities are proportional to $g(E_F)$. This is because it is only the electrons close to E_F that are able to do anything different from the ground state at finite T , (when $T \ll T_F$) and $g(E_F)$ gives the density of such electrons.

For the free electron gas we had

$$g(E) = \frac{3}{2} \frac{m}{E_F} \sqrt{\frac{E}{E_F}}$$

$$\text{so } g(E_F) = \frac{3}{2} \frac{m}{E_F}$$

$$\Rightarrow C_V \sim 3m \frac{k_B^2 T}{E_F} = 3m k_B \left(\frac{T}{T_F} \right)$$

Compare this to a classical ideal gas where

$$C_V = \frac{3}{2} m k_B$$

We see that the specific heat of the quantum fermi gas is $\sim (T/T_F) \sim 10^{-2}$ times smaller at room temperature.

Also, whereas the classical C_V is independent of temperature, the quantum C_V is linear in temperature. This difference will help fix some of the problems with the classical Drude model.

A more precise calculation using the Sommerfeld expansion (see Ashcroft + Mermin pgs 45-47, Eq (2.81) or Kittel pgs 142-144) gives the result

$$C_V = \frac{\pi^2}{3} g(E_F) k_B^2 T$$

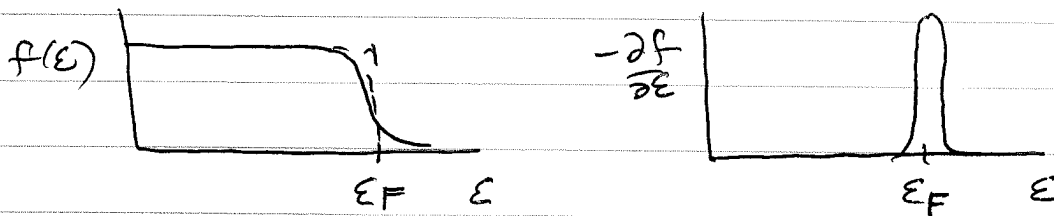
Quick derivation (from Kittel)

$$\textcircled{1} \quad u = \int_0^{\infty} d\varepsilon g(\varepsilon) \varepsilon f(\varepsilon) \Rightarrow \frac{du}{dT} = \int_0^{\infty} d\varepsilon g(\varepsilon) \varepsilon \frac{\partial f}{\partial T}$$

$$\textcircled{2} \quad m = \int_0^{\infty} d\varepsilon g(\varepsilon) f(\varepsilon) \Rightarrow 0 = \frac{\partial m}{\partial T} = \int_0^{\infty} d\varepsilon g(\varepsilon) \frac{\partial f}{\partial T}$$

Subtract E_F times $\textcircled{2}$ from $\textcircled{1}$

$$C_V = \frac{du}{dT} = \frac{du}{dT} - 0 = \int_0^{\infty} d\varepsilon g(\varepsilon) (\varepsilon - E_F) \frac{\partial f}{\partial T}$$



$\frac{\partial f}{\partial T} \approx 0$ except for energies within $k_B T$ of E_F

\Rightarrow approx $g(\varepsilon)$ in integral by constant $g(E_F)$

$$\frac{\partial u}{\partial T} = g(E_F) \int_0^{\infty} d\varepsilon (\varepsilon - E_F) \left(\frac{\partial f}{\partial T} \right)$$

$$\frac{\partial f}{\partial T} = \frac{\partial}{\partial T} \left[\frac{1}{e^{(\varepsilon - \mu)/k_B T} + 1} \right] \approx \left(\frac{\partial f}{\partial x} \right) \left(\frac{\varepsilon - \mu}{k_B T^2} \right)$$

where $x = \left(\frac{\varepsilon - \mu}{k_B T} \right)$. ignore $\mu(T)$ dependence on T
since $\mu \approx E_F$, corrections
give higher order terms

At low $T \ll T_F$, we can replace $\mu(T) \approx E_F$

$$\begin{aligned} \text{So } \frac{\partial u}{\partial T} &= g(E_F) \int_{-\infty}^{\infty} d\varepsilon \frac{(\varepsilon - E_F)^2}{k_B T^2} \left(\frac{-\partial f}{\partial x} \right) \\ &= g(E_F) k_B^2 T \int_{-\infty}^{\infty} dx x^2 \left(\frac{-\partial f}{\partial x} \right) \\ &= g(E_F) k_B^2 T \underbrace{\int_{-\infty}^{\infty} dx x^2 \frac{e^x}{(e^x + 1)^2}}_{= \frac{\pi^2}{3}} \end{aligned}$$

$$\Rightarrow C_V = \frac{\partial u}{\partial T} = \frac{\pi^2}{3} g(E_F) k_B^2 T$$

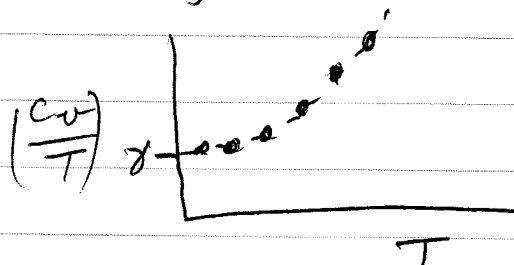
Above is just the contribution to the specific heat of a metal from the conduction electrons. The total specific heat of the metal also includes contributions from vibrations of the ions. This ionic contribution goes like $3n_{\text{ion}}k_B$ at high temperatures ($n_{\text{ion}} = n/z$ is ion density; $z =$ electron valence) but at low T , we will find that due to quantum effects it goes like T^3 . So at high $T \sim$ room temp, the ionic contribution to metallic specific heat dominates the electronic part, and we cannot measure the electronic part to test it against theory. But at low T , one can measure C_V and fit it to the form

$$C_V = \gamma T + AT^3$$

function of
density, n

where our calculation predicts $\gamma = \frac{\pi^2}{2} \frac{nk_B}{T_F}$

one plots $\frac{C_V}{T} = \gamma + AT^2$ and determines γ by extrapolating to the $T \rightarrow 0$ intercept



one finds (see A+M table 3.3)

metal	γ (theory)	γ (expt)	$(10^{-4} \frac{\text{cal}}{\text{mole } ^\circ\text{K}^2})$
Li	1.8	4.2	
Na	2.6	3.5	
K	4.0	4.7	
Rb	4.6	5.8	
Mn	1.5	40.	
Fe	1.5	12	
Bi	4.3	0.2	

The alkali metals (Li, Na) come out reasonable
 The noble metals (Cu, Ag, Au) come out reasonable
 Both alkali & noble metals have valence $Z=1$

But there are much bigger discrepancies for

Fe, Mn, Bi, Sb

\uparrow ~ 10 times too big \uparrow ~ 10 times too small

The difference between γ_{theory} and γ_{expt} is usually due to the difference between the real $g(E_F)$ and the free electron approximation for $g(E_F)$. This is because the potential energy due to the ions shifts the energy spectrum E_k from the simple free electron $\hbar^2 k^2 / 2m$. We can often approx the true E_k as $\frac{\text{near } E_F}{\text{as}} \hbar^2 k^2 / 2m^*$ where m^* is the effective mass that can be either bigger or smaller than the free electron mass.

Transport Properties of the Sommerfeld Model

We have seen that we must use quantum Fermi-Dirac statistics to describe the thermodynamic behavior of conduction electrons.

What about the equation of motion? Semiclassical arguments can be made to show that in many cases it remains OK to use the classical Drude equation of motion. We need to be able to construct electron wave packets which are localized on the desired spatial length scales.

We saw that the typical electron energy is set by the Fermi energy E_F . Hence the typical momentum is $p \sim p_F$. If we make a wave packet with $\Delta p \ll p_F$ then spread in spatial position is

$$\Delta x \sim \frac{\hbar}{\Delta p} \gg \frac{\hbar}{p_F} \sim r_s$$

So the electron cannot be localized to atomic length scales, but can be localized on the length scales of macroscopically varying electric, magnetic fields or temperature gradients $\sim 1000 \text{ \AA}$. So OK for motion in EM waves in visible spectrum

but not for X-rays ($\lambda_{\text{xray}} \sim \text{\AA}$).

We also need $\Delta x \ll l$ the mean free path.
Classical motion may fail when $l \sim 10 \text{\AA}$.

$$l \approx v \tau \quad \text{where } \tau \sim 10^{-14} \text{ sec and}$$

$$v \sim v_F \sim 10^8 \text{ cm/sec}$$

$$\text{so } l \sim 10^8 \cdot 10^{-14} \text{ cm} = 10^{-6} \text{ cm} = 100 \text{\AA}$$

so looks OK.

Using Drude ~~of~~ Equation motion + Fermi Dirac Statistics we then have

① dc and ac electric conductivity same as for classical Drude model, since no thermo was involved

② For thermal conductivity we had

$$\kappa = \frac{1}{3} v^2 \tau c_v$$

Now we should use $v \sim v_F$, $v_F^2 = \frac{2 E_F}{m}$

$$\text{and } c_v = \frac{\pi^2}{3} \left(\frac{k_B T}{E_F} \right) m k_B$$

$$\Rightarrow \kappa = \frac{1}{3} \frac{2 E_F}{m} \frac{\pi^2}{2} \left(\frac{k_B T}{E_F} \right) m k_B \tau$$

$$= \frac{\pi^2}{3m} \tau m k_B^2 T$$

and Wiedemann-Franz coefficient is

$$\frac{\kappa}{\sigma T} = \frac{\frac{\pi^2}{3m} \tau m k_B^2}{\frac{n e^2 \tau}{m}} = \frac{\pi^2}{3} \left(\frac{k_B}{e} \right)^2 = 2.44 \times 10^{-8} \text{ watt}^2 / \text{ohm} / \text{K}^2$$

excellent agreement
with experiment

③ Thermopower

$$Q = -\frac{C_V}{3me} = -\frac{\pi^2}{6} \frac{k_B}{e} \left(\frac{k_B T}{E_F} \right)$$
$$= -1.42 \left(\frac{k_B T}{E_F} \right) \times 10^{-4} \text{ volt / K}$$

more reasonable result than
classical value