Theory of thermoelectric phenomena

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Summary of the scaling results for the CF split Anderson model

In the high-T phase, all the CF states are occupied, the f-state is fully degenerate.

- In Ce compounds, $J = 5/2$ and $N = 6$.
- In Yb compounds, $J = 7/2$ and $N = 8$.

The CF splitting defines the temperature scale $k_B T_\Delta = \Delta/2$.

- For $T \geq T_\Delta$, the scaling trajectory is defined by the high-T Kondo scale, $T_K^H \ll \Delta$.
- Below $T_\Delta$, the excited CF states depopulate, the degeneracy of f-states drops from $N$ to, say, $N_m = 2$.
- For $T_K < T < T_\Delta$, the Kondo physics is governed by $T_K \ll T_K^H$.
- At $T \approx T_K$, the system makes a transition to the low-T phase, properties are governed by the low-T Kondo scale $T_K \ll T_K^H$. 
Summary of the scaling results for the Anderson model

For $T \gg T_K$, the f-electrons are excluded from the Femi volume. The current is carried by conduction electrons which are exchange-scattered on localized f-states. The system can be described by the Kondo lattice.

For $T \leq T_K$, the large entropy of degenerate f-states is removed by the crossover into the coherent Fermi liquid regime.

In the Fermi liquid phase, $T \ll T_K$, the f-electrons are included in the Fermi volume and the current is carried by hybridized quasiparticles. The resistivity is caused by the quasiparticle-quasiparticle scattering.
However, the high-temperature entropy might be removed in many different ways, not just by Kondo effect, as the system might have several competing ground states.

If $T_K$ is small, the crossover in the FL state might be intercepted by a phase transition, like the superconducting, magnetic or structural one.

A small change in parameters can change in the ground state. Thus, these systems are candidates for quantum critical behavior.

Pressure can change the degeneracy of the $f$ states. The CF splitting can be quenched above some critical pressure. Thus, temperature dependence of transport coefficients is modified by pressure.
Low-temperature properties of Kondo systems

In the coherent regime, the valence fluctuators and heavy fermions follow the Fermi liquid laws, typical of Fermi liquids.

- Enhanced Pauli susceptibility:
  
  \[ \chi_0 = \left( \frac{g \mu_{\text{eff}}}{T_0} \right)^2 \]

- Linear specific heat:
  
  \[ C_v = \gamma T, \quad \text{where} \quad \gamma = \frac{1}{T_0} \]

- Parabolic resistivity:
  
  \[ \rho(T) = \rho_0 + AT^2 \]

- Linear thermopower:
  
  \[ \alpha(T) = \alpha_0 T, \quad \text{where} \quad \alpha_0 = \frac{1}{T_0} \]

All the properties are controlled by a single parameter \( T_0 \).
Consequences of the Fermi liquid laws

In the limit $T \to 0$, the ratios of various physical quantities exhibit universal behavior

- **Wiedemann-Franz law:**
  
  $$\mathcal{L}_0 = \frac{\kappa}{\sigma} = \frac{\pi^2 k_B^2}{3e^2} T$$

- **Wilson ratio:**
  
  $$R_W = \chi_0 / \gamma$$

- **Kadowaki-Woods ratio:**
  
  $$R_{KW} = \rho(T) / (\gamma T)^2$$

- **q-ratio:**
  
  $$q = |e| \lim_{T \to 0} \alpha_0 / \gamma$$
Figure: (A) Coefficient of the $T^2$ term in the resistivity plotted as a function of $\gamma$, for heavy fermion systems with various degeneracies of the $f$-level. The continuous lines represent the predictions of the Kadowaki–Woods relation and its generalization to $N > 2$. (B) Thermoelectric power divided by temperature (close to $T = 0$ K), plotted as a function of $\gamma$, for a number of elemental metals, heavy fermion systems, metallic oxides and organic superconductors.
Theory should account for the FL laws and explain the deviations from the universal ratios.

Another puzzling feature of the Fermi liquid state is the pressure dependence shown by the residual resistance $\rho_0(p)$ and the coefficient of the $T^2$ term in the low-temperature resistivity

$$A(p) = \frac{\rho(p) - \rho_0(p)}{T^2}$$

The pressure data exhibit anomalies at the crossover between various regimes.

The properties of the FL state and the high-temperature properties seem to be correlated.
Figure: Enhancement of residual resistivity in several samples, scaled to a universal pressure dependence. \((a\text{ and } \rho_0^* \text{ are normalizing factors.})\)
Pressure dependence of $T^2$ coefficient in CePd$_2$Ge$_2$

FIG. 4. Temperature coefficient $\tilde{A}$ obtained from a fit, $\rho_{\text{mag}}(T) = \rho_0 + \tilde{A} T^n$, to the data of CePd$_{2+x}$Ge$_{2-x}$ ($x=0$ and 0.02) below 2 K vs pressure $p$. $\tilde{A}$ peaks at 11.0 GPa and 13.8 GPa for $x=0.02$ and $x=0$, respectively. Inset shows the exponent $n$ used to describe $\rho_{\text{mag}}(T)$ at different pressures.
Figure: Pressure dependence of the residual resistivity $\rho_0(p)$ and the coefficient $A(p)$ in the relation $\rho = \rho_0 + AT^2$ for CeCu$_2$Ge$_2$. $\rho_0$ has a maximum where $A(p)$ has a discontinuity.
Correlation between the FL and the high-T scales

Figure: Coefficient $A(\rho)$ plotted versus the temperature $T_\rho(\rho)$ of the first resistivity maximum for CeCu$_2$Si$_2$ and CeCu$_2$Ge$_2$ at various pressures. The straight lines indicate a scaling $A \propto T_\rho^{-2}$. Pressure increases toward the right-hand side of the scale (high $T_\rho$).
To describe the HF and the FV systems in the coherent regime we use the periodic Anderson model with SU(N) symmetry. $N$ is the effective degeneracy of the 4f states.

The hybridization can mix only the states which belong to the same irep, so the model describes $N$ conduction channels connected in parallel.

For $n_f$ electrons per Ce ion and $n_c$ electrons per conduction channel, there are $n = Nn_c + n_f \leq 2$ electrons at each lattice site.

The description based on this model is a drastic simplification with respect to the real materials and the model applies only to the low-energy part of the excitation spectrum.

However, the model describes the effects of local correlations and if these correlations are causing the low-temperature anomalies, the qualitative features should be captured by the solution.
The difficulty is that, unlike the single impurity models, the lattice models with on-site correlation can be solved accurately only in some limiting cases.

Thus, we do not really know if these effective models really explain the anomalies that one finds in the experimental data.

Recently, there are serious efforts to combine the band structure of the VF and HF compounds with the many body treatment of local correlations, the so called DMFT approach.

The DMFT provides the solution of the lattice problem with on-site correlations at all temperatures and is very promising. But more work is to be done.
The solution at low temperatures can be obtained using Landau’s quasiparticle description.

The universal behavior shown by the low-temperature experimental data points to a single relevant Fermi liquid scale. Choose,

$$T_0 \simeq \frac{1}{\gamma}$$

The specific heat coefficient $\gamma$ can either be taken from the experiment or calculated, together with the chemical potential $\mu$, for the system in thermal equilibrium.

Once $T_0$ is fixed, we can express all other physical quantities as the universal functions of $T/T_0$ or $\omega/T_0$. 
The equilibrium properties are relatively easy to calculate. Transport is more complicated.

The transport properties depend on

- the renormalized density of conduction states $N_c(\omega)$,
- the transport relaxation time $\tau(\omega, T)$,
- the average velocity of current carrying states $\langle v_{kF} \rangle$.

The FL theory relates all these quantities to $T_0$ and the unrenormalized density of conduction states, that we know,

$$N_c^0(\omega) = \frac{1}{N_i V_{\text{cell}}} \sum_k \delta(\omega - \epsilon_k)$$

The Fermi surface in each channel is fixed by the number of particles $n_c + n_f / N$. 
Calculation of transport coefficients by the DMFT

\[ L_{mn} = e^{4-m-n} \int d\omega \left( -\frac{\partial f}{\partial \omega} \right) \omega^{m+n-2} \Lambda_{\text{tr}}(\omega, T) \]

For 3-d systems the computations of \( \Lambda_{\text{tr}}(\omega, T) \) are impossible and we use dynamical mean field approximation which neglects all vertex corrections. In this approximation, the \( \omega, T \to 0 \) limit yields

\[ \Lambda_{\text{tr}}(\omega, T) = \frac{1}{3} v_F^2 N_c(\omega) \tau(\omega, T), \]

The transport relaxation time is defined by the self energy of conduction electrons

\[ \tau(\omega, T) = \frac{-1}{\text{Im} \, \Sigma_c(\omega^+, T)} \]

In DMFT, \( \text{Im} \, \Sigma_c(\omega^+, T) \) is momentum-independent.

To get the renormalized density of states we need the Green’s function

\[ N_c(\omega) = -\frac{1}{\pi N_i V_{\text{cell}}} \sum_k \text{Im} \, G_c(k, \omega^+). \]
The Green’s functions are defined by their equations of motion. Conduction electrons Green’s function of PAM is given by,

\[ G_c(k, z) = \frac{1}{z - \epsilon_k + \mu - \Sigma_c(z)}, \]

The self energy is

\[ \Sigma_c(z) = \frac{V^2}{z - E_f + \mu - \Sigma_f(z)}. \]

The f-electron Green’s function is

\[ G_f(k, z) = \frac{1}{z - E_f - \Sigma_f(z) + \mu - V^2/(z - \epsilon_k + \mu)}, \]

\( \Sigma_f(z) \) is k-independent for local hybridization.

The problem is to find \( \Sigma_f(z) \). In general, this problem can’t be solved.
The Fermi liquid approach

The low-energy excitations of the system are defined by the singularities of $G_c(k, z)$ and $G_f(k, z)$.

At low temperatures, the singularities can be obtained in the quasiparticle approximation, which neglects the decay of QP and assumes

$$\text{Im } \Sigma_f(\omega) \approx 0$$

Expanding $\text{Re } \Sigma_f(\omega)$ around $\omega = 0$ yields

$$\omega - [E_f + \text{Re } \Sigma_f(\omega) - \mu] \approx (\omega - \tilde{\omega}_f) Z_f^{-1} + O(\omega^2)$$

The quasiparticle weight (the mass enhancement) is

$$Z_f^{-1} = 1 - \partial \Sigma_f / \partial \omega |_{\omega=0}$$

The shift of energy levels due to the interaction is

$$\tilde{\omega}_f = [E_f + \text{Re } \Sigma_f(0) - \mu] Z_f$$
The Fermi liquid approach

The singularities of $G_c(k, z)$ and $G_f(k, z)$ can now be obtained from the solution of the secular equation

$$(z - \epsilon_k + \mu)(z - E_f - \Sigma_f(z) + \mu) - V^2 = 0$$

In the quasiparticle approximation, the poles of both Green’s functions are determined by

$$(\omega - \epsilon_k + \mu)(\omega - \tilde{\omega}_f) - \tilde{V}^2 = (\omega - \Omega^-)(\omega - \Omega^+ ) = 0.$$ 

The renormalized hybridization is $\tilde{V} = V \sqrt{Z_f} \ll V$.

The quasiparticle energies are given by

$$\Omega^\pm_k = \frac{1}{2} \left[ (\epsilon_k - \mu + \tilde{\omega}_f) \pm \sqrt{(\epsilon_k - \mu - \tilde{\omega}_f)^2 + 4 \tilde{V}^2} \right]$$

Quasiparticle spectrum has two branches separated by an indirect gap

$$\Delta_{QP} = \frac{\tilde{V}^2}{\mu + \tilde{\omega}_f}$$
For a system with an electron-hole symmetry, $\mu$ is in the gap.

Away from the e-h symmetry, the Fermi surface of QP is given by

$$\Omega_{k_F}^- = 0 \quad \text{or} \quad \Omega_{k_F}^+ = 0$$

Close to $k_F$, the QP dispersion is very weak.

In the QP approximation, the low-energy excitations are represented by the non-interacting Fermions with an effective bandwidth $\tilde{V}_f^2/W$.

In Cerium, $n_f < 1$, $E_f + \text{Re} \Sigma_f(0) < \mu$, and $\mu$ is in the lower QP branch. $k_F$ is close to the zone boundary.

In Ytterbium, $n_f > 1$, $E_f + \text{Re} \Sigma_f(0) > \mu$, and $\mu$ is above the gap. $k_F$ is close to the center of the Brillouin zone.
Quasiparticles are parametrized by $Z_f$ and $\tilde{\omega}_f$, which are not known, since we didn’t calculate $\text{Re} \ \Sigma_f(\omega)$.

But we can calculate the QP specific heat, $C_V = \partial E / \partial T$, assuming that thermally excited QP increase the thermal energy of the system by $\Delta E = (k_B T)^2 N N^{QP}(0)$. Thus,

$$\gamma = \frac{\pi^2 k_B^2}{6} NN^{QP}(0) \equiv \frac{\pi^2 k_B^2}{3\mathcal{V}_\text{cell}} \frac{1}{k_B T_0}$$

Like for any free fermions, all the low-temperature properties of quasiparticles can be expressed in terms of $T_0$.

In the QP approximation,

$$A_c(\mathbf{k}, \omega) \approx a^c_\mathbf{k}(\omega) \delta(\omega - \Omega^{\pm}_\mathbf{k})$$

and

$$A_f(\mathbf{k}, \omega) \approx a^f_\mathbf{k}(\omega) \delta(\omega - \Omega^{\pm}_\mathbf{k}),$$

This yields

$$\frac{1}{k_B T_0} \approx \frac{N}{2} \frac{\mathcal{V}_\text{cell} N_f(0)}{Z_f}$$
The renormalized DOS of conduction electrons is

\[ N_c(\omega) = \frac{1}{N_i V_{\text{cell}}} \sum_k A_c(k, \omega) \]

This gives

\[ N_c(\omega) = N^0_c \left( \omega + \mu - \frac{\tilde{V}^2}{\omega - \tilde{\omega}_f} \right) \]

The renormalized c-DOS is given by the DOS of the unhybridized conduction states at shifted frequency. \( N^0_c(\omega) \) is small everywhere.

Similarly,

\[ N_f(\omega) = \frac{Z_f \tilde{V}^2}{[\omega - \tilde{\omega}_f]^2} N_c(\omega). \]

and

\[ \frac{N_f(0)}{Z_f} = \frac{1}{\tilde{\omega}_f} \frac{\tilde{V}^2}{\tilde{\omega}_f} N^0_c(\mu + \frac{\tilde{V}^2}{\tilde{\omega}_f}) \]

The shift \( \tilde{V}^2/\tilde{\omega}_f \) is easy to calculate and has a simple interpretation.
Chemical potential $\mu$ is fixed by the total number of particles, so that

$$n = N n_c + n_f = N V_{\text{cell}} \int_{-\infty}^{\mu} d\omega \left[ N_c(\omega) + N_f(\omega) \right].$$

The renormalized Fermi surface is given by the solution of the equation

$$\epsilon_{k_F} = \mu + \frac{\tilde{V}^2}{\tilde{\omega}_f} = \mu_L$$

By Luttinger theorem, the number of k-points enclosed by $k_F$ equals $n$. Since $\epsilon_k$ is the non-interacting dispersion, we can find $\mu_L$ using

$$n = N V_{\text{cell}} \int_{-\infty}^{\mu_L} d\omega N^0_c(\omega).$$

Thus, $\tilde{V}^2/\tilde{\omega}_f$ is the shift of the chemical potential due to hybridization,

$$\Delta \mu = \mu_L - \mu = \frac{\tilde{V}^2}{\tilde{\omega}_f}$$

Since $n$ and $N^0_c(\omega)$ are known, $\Delta \mu$ is easy to find, and we can write

$$2\tilde{\omega}_f = \Delta \mu V_{\text{cell}} N^0_c(\mu_L) T_0.$$
For transport properties, we can’t ignore the quasiparticle damping due to the quasiparticle-quasiparticle scattering.

In infinite dimensions, the perturbation theory for $\Sigma_f$ gives the lifetime

$$-\text{Im} \Sigma_f (\omega, T) \approx \frac{\pi}{2} \left[ \omega^2 + (\pi k_B T)^2 \right] (N - 1) [\nu_{\text{cell}} N_f (0)]^3 \Gamma_f^2$$

$\Gamma_f$ is the four-point scattering vertex due to Coulomb interaction.

The Ward identity relates $\Gamma_f$ and $Z_f^{-1}$,

$$Z_f^{-1} = (N - 1) \nu_{\text{cell}} N_f (0) \Gamma_f.$$

Using $N_f (0) / Z_f \sim 1 / T_0$ obtains

$$\frac{1}{k_B T_0} \approx \frac{N(N - 1)}{2} [\nu_{\text{cell}} N_f (0)]^2 \Gamma_f.$$

So we can relate the QP damping to $T_0$. 
The transport relaxation time, at low frequencies, is given by

$$\tau(\omega, T) \simeq \tau_0(T) \left(1 - \frac{\omega}{\tilde{\omega}_f}\right)^2 \left(1 - \frac{\omega^2}{\pi^2 k_B^2 T^2}\right),$$

where

$$\tau_0(T) = \frac{\hbar(N - 1)N^2\mathcal{V}_{cell}\mathcal{N}_c^0(\mu_L)}{2\pi^3} \frac{T_0^2}{T^2}.$$

Since we know

$$\mathcal{N}_c(\omega) = \frac{(\omega - \tilde{\omega}_f)^2}{Z_f \tilde{V}^2} \mathcal{N}_f(\omega),$$

the transport coefficients can be calculated by Sommerfeld expansion.

The Fermi surface average of the unrenormalized velocity squared is

$$\langle v_{k_F}^2 \rangle = \int d^d k \ \delta(k - k_F) \ v_k^2,$$

This is easy to calculate for any dispersion.

Thus, we have all that is needed for transport coefficients in the Fermi liquid regime.
The resistivity of $N$ parallel conducting channels is

$$
\rho(T) = \frac{9\pi^3 V_{\text{cell}}}{\hbar e^2 v^2_F N(N-1)N\mathcal{N}_c N_c^0(\mu_L)^2} \left( \frac{T}{T_0} \right)^2 = A^* \left( \frac{T}{T_0} \right)^2,
$$

If we tune the model parameters but preserve the Fermi volume of each channel (keeping $N/n$ constant), the Luttinger theorem ensures that $\mu_L$ and $v^2_F$ do not change. In that case, $A^*$ is constant and the coefficient of the $T^2$ term changes owing to variations in $T_0$.

The factor $A^*$ depends on the external parameters.

- additional impurity scattering or lattice expansion can localize the $f$-states and exclude them from the Fermi volume.
- the change in the effective degeneracy of the $f$-states can shift the Fermi surface of the resonant channels.
- The “jump” in the Fermi volume changes $\mu_L$ and $v^2_F$, and impacts the resistivity.
The Sommerfeld expansion for $L_{12}$ and $L_{11}$ gives the Seebeck coefficient

$$\alpha(T) \simeq \pm \frac{4\pi^2 k_B}{|e| N \Delta \mu \mathcal{N}_{\text{cell}} \mathcal{N}_c^0(\mu_L)} \frac{T}{T_0}.$$  

The sign of $\alpha(T)$ is positive for Ce- and Eu-based systems and negative for Yb-based ones.

$\Delta \mu$ has an explicit parameter dependence. Strictly speaking, $\alpha(T)$ is not a universal function of $T/T_0$.

For $n_f \approx 1$ and large $N$ we can write

$$\alpha(T) = \pm \frac{4\pi^2 k_B}{n_f |e|} \frac{T}{T_0}.$$

The closer the system to half-filling, the larger the slope, $\alpha/T \propto 1/T_0$. Large slope does not imply large $\alpha(T)$ at finite temperatures, since the FL laws are only valid for $T \ll T_0$. Close to half-filling, $T_0$ is exponentially small.
The thermal conductivity in the clean limit is
\[ \kappa(T) = T \sigma(T) \mathcal{L}_0(T), \]

with the effective Lorenz number
\[ \mathcal{L}_0(T) = \bar{\mathcal{L}}_0 \left[ 1 - \frac{32\pi^2}{n_f^2} \left( \frac{T}{T_0} \right)^2 \right], \]

and \( \bar{\mathcal{L}}_0 = \frac{\pi^2 k_B^2}{2e^2} \).

The deviation form the Wiedemann–Franz law is due to the QP damping. \( T \to 0 \) limit yields the relation \( \kappa(T) \propto T \sigma(T) \) but the coefficient of the \( T^2 \) term is very large, so there are substantial deviations from the Wiedemann–Franz law much below \( T_0 \).

The deviations from the WF law enhance the figure of merit. Neglecting the phonons in the Fermi liquid regime, we can write
\[ ZT = \frac{\alpha^2(T)}{\mathcal{L}(T)}. \]
The KW ratio, \( \rho(T)/(\gamma T)^2 \), reads

\[
R_{KW} = \frac{81}{h k_B^2 e^2} \frac{\mathcal{V}_{\text{cell}}^3}{v_F^2 [N \mathcal{V}_{\text{cell}} \mathcal{N}_c^0(\mu_L)]^2} \frac{1}{\frac{1}{2} N(N - 1)}
\]

\( R_{KW} \) exhibits an explicit dependence on the ground-state degeneracy and the average velocity (squared), and an implicit dependence on the Fermi volume (i.e., on the carrier concentration \( n \)).

The \( q \)-ratio is

\[
q = \lim_{T \to 0} |e\alpha|/\gamma T \simeq 12/[N \Delta \mu \mathcal{N}_c^0(\mu_L)],
\]

\( \Delta \mu \) has an explicit parameter dependence, so that the \( q \)-ratio can deviate from the universal value.

The transfer of \( f \)-electrons into the conduction band makes the \( q \)-ratio pressure-dependent.

Data on the hydrostatic pressure-dependence of the \( q \)-ratio are not yet available, but the deviations from the universal value are indicated by chemical pressure experiments.
Pressure dependence of the resistivity

In the FL regime, the resistivity of real materials has a contribution due to the QP-QP scattering and the scattering of the QPs by lattice imperfections.

In the AFM phase, there is also a contribution due to the scattering on spin waves.

For \( T \leq T_N \), the f-electrons are excluded from the Fermi volume and parts of the FS are gapped. The exchange coupling increases with pressure and so does \( \rho(T) \) in the AFM phase.

At the critical pressure, the system becomes a paramagnet and the gaps at the FS close. At \( p_c \), the conductivity increases sharply.

In the paramagnetic regime, the experimental data are represented as

\[
\rho(T) = \rho_0 + A(p) T^2
\]

The theoretical expression is

\[
\rho(T) = \rho_0 + A^*(p) \left( \frac{T}{T_0} \right)^2
\]
For $p \geq p_c$, the transport relaxation time of has a contribution due to QP-QP scattering and the residual scattering on lattice imperfections.

For $T \leq T_\Delta$, the QP-QP scattering is described by the SU(N) Anderson model with $N_m$ channels with heavy electrons and $N - N_m$ channels with nearly free conduction states.

Each hybridized channels accommodates $n_c + n_f/N_m$ electrons. The FS is close to the zone boundary and $\mu_L$ close to the band edge. In these channels, $v_F^2$ and $N_c^0(\mu_L)$ are small, and $A(p)$ is large.

$N - N_m$ channels with nearly free-electron states have $T$-independent resistivity. The current is mainly carried by these channels. The small effective mass and high Fermi velocity make $\rho_0$ very small.

$T_0$ is not provided by our calculation but there is theoretical evidence that the FL scale of the coherent regime is about the same as the $T_K$ scale of the incoherent regime that we know how to calculate.
As long as only the lowest CF states are occupied, pressure increases $T_0$ but does not change $A^*$. The pressure dependence of the experimental parameter $A(p) = A^* / T_0^2$ is obtained from $T_0(p) \simeq T_K(p)$.

The resistivity of Kondo impurities has a maximum at about $T_\rho \simeq T_K$, which explains the linear correlation between $A(p)$ and $T_\rho^{-2}$.

For $p_c \leq p \leq p_\nu$, the residual resistivity is due to the scattering of nearly free electrons, in $N - N_m$ channels, on lattice imperfections. Thus, $\rho_0$ is nearly constant.
At $p \simeq p_v$, the CF excitations collapse and f-electrons are evenly distributed over $N = N_m + N_M$ channels. That is, $N - N_m$ channels acquire a heavy mass, so that their Fermi surface expands, and $\nu_F$ drops. Thus, $\rho_0$ increases dramatically at $p_v$.

At $p_v$, the coefficient $A^*$ is reduced, because the hybridized channels accommodate now $n + n_f/N$ electrons, rather than $n + n_f/N_m$ electrons. The Fermi volume shrinks, so that $\nu_F^2$ and $\mathcal{N}_c^0(\mu_L)$ increase with respect to their values at $p < p_v$.

The coefficient $A(p) = A^*(p)/T_0^2$ is even more drastically reduced at $p_v$, because $T_0$ is now equal to $T_K^H \gg T_K$.

A further increase in pressure reduces $A(p)$ but does not change $A^*$, since $\mu_L$ is fixed by the Luttinger theorem for $n_c + n_f/N$ states per channel. This explains the linear correlation between $A(p)$ and $(T_\rho)^{-2}$ found for $p \geq p_v$.

The reduction of the QP mass ($\propto 1/T_0$) explains the rapid decrease of $\rho_0$ for $p > p_v$.
The transition is shown in detail in the inset. Fig. 3. The superconducting transition of CeCu$_2$Si$_2$ and CeCu$_2$Ge$_2$ close to the main transition in zero field is located at different values, as shown in Fig. 3 for two ordering. Without any other results, the change of slope of below 1.8 K the superconductivity sets in. The ohmic response of sample & drop is not due to superconducting residual resistivity. Another example is the occurrence of this peak can be associated to the normal phase resistivity, a drop of $\rho_0$ shows a standard transition at 1.4 K, its resistivity has dropped below our threshold. The occurrence of this peak can be associated to the main transition. For example, CeNi$_2$Ge$_2$ can also exhibit different behavior. Below 0.55 K, where $\rho_0$ varies increases only for YbCu$_2$Si$_2$. For a real crystal, the residual resistivity $\rho_0$ (1/\(\Omega \cdot \text{cm}\)) changes. It appears that the superconductivity is likely different when the superconductivity sets in. The occurrence of this peak can be associated to the main transition. For high enhancement seems. For high field this contribution shifts to lower $P_P_c$ (GPa). The occurrence of this peak can be associated to the main transition. For high enhancement seems. For high field this contribution shifts to lower $P_P_c$ (GPa).
In CeRu$_2$Ge$_2$ and CePd$_{2-x}$Ge$_{2-x}$, the transition at $p_c$ is directly into the fully degenerate 4f spin-orbit multiplet. At $p_c$, the conduction electrons become heavy, so $A(p)$ and $\rho_0(p)$ become very large.

Above $p_c$, the coefficient $A^*$ is pressure independent, so that the rapid decrease of $A(p)$ and $\rho_0(p)$ is caused by the increase of $T_0 \simeq T_K$.

YbCu$_2$Si$_2$, at ambient pressure, is a valence fluctuator. An increase of pressure stabilizes 4f$^{13}$-configuration and gives rise to a HF state in all 8 channels of the $j = 7/2$ multiplet. Pressure increases the effective mass, so that $\rho_0$ increases with pressure.

At higher pressure, the CF excitations start to appear, and, eventually, the effective degeneracy of the 4f-hole drops from $N = 8$ to $N = 2$. The transport is through the $N - 2$ “nonresonant” sub-bands, so that $\rho_0$ drops abruptly.