

Physics 745: Special Topics in Theoretical Physics

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## Paper on Fermi liquids

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# Contents

<b>1</b>	<b>Introduction &amp; the Quasiparticle Concept</b>	<b>2</b>
1.1	Non-interacting Fermi Gas . . . . .	2
1.2	Adiabatic Continuity . . . . .	3
<b>2</b>	<b>Why Quasiparticle survive</b>	<b>5</b>
<b>3</b>	<b>The Landau Energy Functional</b>	<b>7</b>
3.1	The Landau Parameters . . . . .	8
<b>4</b>	<b>Predictions &amp; Observables</b>	<b>9</b>
<b>5</b>	<b>Conclusion</b>	<b>9</b>

# 1 Introduction & the Quasiparticle Concept

In this paper, we want to motivate the question of Why Fermi liquids matter. We will want to explain what a Fermi Liquid is, and how it differs from a non-interacting Fermi Gas. We want to build the model behind the Fermi liquid and show that strong electron-electron interactions do not destroy the free-electron model. Lastly, we will find concrete, measurable predictions of specific heat, susceptibility and resistivity using the Fermi liquid model and see if they match with experiments. Before we begin answering these core questions, we look a different question, in order to motivate the idea and to see why we need Fermi liquids. We ask, if interacting electrons are expected to behave chaotically, why do we see that electrons in metals actually act as if they are almost free. We will see that Landau's Fermi Liquid theory explains this remarkable stability. As a note, in this paper, we use  $\mu = E_F$  which we identify at temperature 0 the chemical potential is the Fermi energy.

Classical theory predicts that in a typical metal, electrons interact via strong Columb repulsion comparable to their kinetic energy:

$$E_C = \frac{e^2}{4\pi\epsilon_0 r} = \frac{e^2}{a_0 r_s} \propto \frac{1}{r_s}$$

Comparable to the kinetic fermi energy:

$$E_F = \frac{\hbar^2 k_F^2}{2m} = \frac{\hbar^2}{2ma_0^2 r_s^2} \propto \frac{1}{r_s^2}$$

We can see that the ratio  $\frac{E_C}{E_F} \propto r_s$ . This Columb repulsion, is clearly not a small perturbation in our system. But when we look at experimental data, we get free electron properties. Properties that are identical to those predicted by the non-interacting Fermi Gas model. We find specific heat as a linear dependence on temperature,  $C_V \propto T$ , a constant magnetic susceptibility,  $\chi$ . And a well-defined Fermi surface.

The Fermi Liquid defined as an interacting fermionic system whose low-energy excitations are adiabatically connected to those of a free Fermi gas. To see that this definition holds, lets first define and model both the non-interacting Fermi Gas model as well as detail what the Fermi surface is.

## 1.1 Non-interacting Fermi Gas

Consider a system of  $N$  non-interacting fermions. The grand canonical Hamiltonian is purely kinetic (Coleman eqn 4.43):

$$H = H_S - \mu N = \sum_{k,\sigma} (E_k - \mu) c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma}$$

where  $c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma}$ , is the creation operator and annihilation operator respectively, and  $\hat{n}_{\mathbf{k}\sigma} = c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma}$  is the number operator. Which counts the number of fermions in the state  $(\mathbf{k}, \sigma)$ . The number of fermions in a state with momentum  $\mathbf{p} = \hbar\mathbf{k}$  is

$$n_k = f(E_k - \mu) \quad f(x) = \frac{1}{e^{\beta x} + 1}$$

during thermal equilibrium. At low temperatures, the Fermi-Dirac function  $f(x)$  represents a step with a jump in occupancy spread over an energy range of order  $k_B T$  around the chemical potential. At  $T = 0$  K, the occupancy of each state is given by:

$$n_k = \theta(\mu - E_k)$$

$\theta(x)$  is the Heaviside step function. As Coleman describes it, it is a step function with an abrupt change in occupation when  $\epsilon = \mu$ , which corresponds to the fact that states with  $E_k < \mu$  are completely occupied and states above this energy are empty. We refer to this zero temperature value of the chemical potential as the *Fermi Energy*. In momentum space,  $k_F$  is the *Fermi Momentum*.

The ground state at zero temperature,  $|\Psi_0\rangle$ , is constructed by filling all momentum states up to the Fermi momentum  $k_F$ :

$$|\Psi_0\rangle = \prod_{|\mathbf{k}| < k_F, \sigma} c_{\mathbf{k}\sigma}^\dagger |0\rangle$$

Excitations above this ground state are produced by addition of particles at energies above the Fermi momentum, or the creation of holes beneath the Fermi momentum. To formally describe these excitations, we have a particle-hole transformation:

$$a_{\mathbf{k}\sigma}^\dagger = \begin{cases} c_{\mathbf{k}\sigma}^\dagger & (k > k_F) \quad \text{particle} \\ \text{sgn}(\sigma) c_{-\mathbf{k}-\sigma} & (k < k_F). \quad \text{hole} \end{cases} \quad (4.49)$$

We can also say that  $b_{k,\sigma}^\dagger = c_{k\sigma}$  for  $k < k_F$  is for hole creation.

We note that, to create a hole with momentum  $\mathbf{k}$  and spin  $\sigma$ , we must destroy a fermion with momentum  $-\mathbf{k}$  and spin  $-\sigma$ . For this Fermi gas, we can continue to define its properties like ground-state density or pressure and etc. as outlined in Coleman. For all of these key features of the Fermi gas, we might expect them all to fall off when interactions become present, but we will see that from the Fermi liquid theory for interacting fermions, we are able to see the key features of this degenerate Fermi gas does survive. We see that instead of the Fermi surface being blurred by particle-particle interactions, we actually have that the Fermi surface is stable (for  $d > 1$ ). The reason being is that electrons at the Fermi surface have no phase space for scattering. Which we will see with quasiparticles surviving. But first, let's discuss Adiabatic Continuity.

## 1.2 Adiabatic Continuity

Having defined the non-interacting limit, we now turn to Landau's derivation of the Fermi Liquid. The central problem is to understand how the eigenstates of the free Hamiltonian  $H_0$  evolve when the interaction  $V$  is introduced. Landau postulated that the interacting state is not a new phase of matter, but rather an adiabatically connected extension of the free gas. Consider a time-dependent Hamiltonian  $H(t) = H_0 + \lambda(t)V$ , where the coupling constant  $\lambda(t)$  increases slowly from 0 at  $t = -\infty$  to 1 at  $t = 0$ . The evolution of the system is governed by the time-evolution operator  $U(t, -\infty)$ . If the turn-on process is sufficiently slow (adiabatic) and no phase transition occurs, the non-interacting ground state  $|\Psi_0\rangle$  evolves smoothly into the interacting ground state  $|\Phi_0\rangle$ :

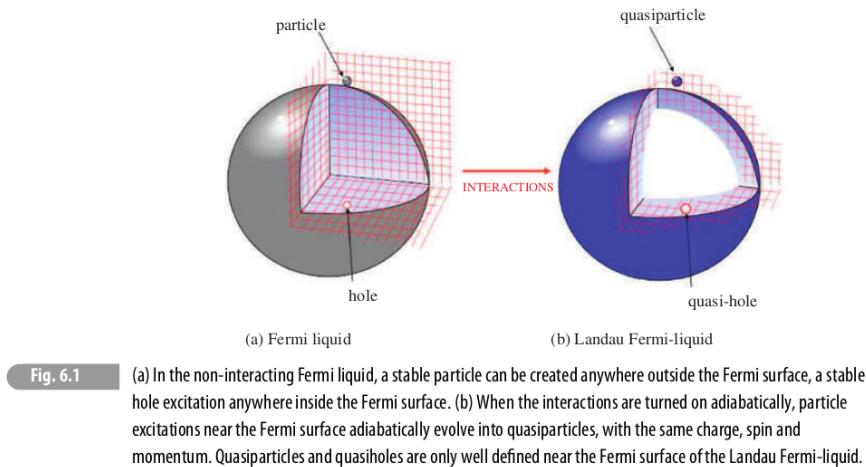
$$|\Phi_0\rangle = U(0, -\infty)|\Psi_0\rangle$$

We see that this mapping extends to the excited states. Consider a single particle excitation of the free gas created by adding a fermion with momentum  $k > k_F$ . This state  $c_{\mathbf{k}\sigma}^\dagger |\Psi_0\rangle$  evolves into a long-lived excitation of the interacting system, which we define as a **quasiparticle**. We can define the quasiparticle creation operator  $a_{\mathbf{k}\sigma}^\dagger$  by transforming the bare creation operator under the same unitary evolution:

$$a_{\mathbf{k}\sigma}^\dagger = U(0, -\infty) c_{\mathbf{k}\sigma}^\dagger U^\dagger(0, -\infty)$$

This transformation yields a state  $|\Psi_{\mathbf{k}\sigma}\rangle = a_{\mathbf{k}\sigma}^\dagger |\Phi_0\rangle$  that corresponds one-to-one with the free electron state. Because the evolution is unitary and preserves the symmetries of the system (translational and rotational invariance), the quasiparticle retains the exact quantum numbers of the bare electron: it carries charge  $e$ , spin  $1/2$ , and momentum  $\mathbf{k}$ .

This structure is illustrated schematically in the figure below. In the non-interacting Fermi gas, particle and hole excitations are well defined throughout momentum space. When interactions are turned on adiabatically, only excitations near the Fermi surface remain sharply defined, evolving into quasiparticles and quasiholes with the same quantum numbers. States far from the Fermi surface lose their single-particle character.



However, the quasiparticle is not identical to a bare electron. Physically, it represents a bare electron dressed by a cloud of virtual particle-hole excitations. This dressing leads to the renormalization of dynamical properties such as the **Effective Mass**. The quasiparticle moves with a renormalized mass  $m^* \neq m$ , which accounts for the inertia of the interaction cloud. At the Fermi surface there is also the **Wavefunction renormalization or spectral weight ( $Z$ )** which can be thought of as the overlap with the state formed by adding a bare particle to the ground state. The spectral weight  $Z$  is defined as:

$$Z = |\langle \Phi_0 | c_{\mathbf{k}\sigma} a_{\mathbf{k}\sigma}^\dagger | \Phi_0 \rangle|^2 < 1$$

In a Fermi gas,  $Z = 1$ . In a Fermi liquid,  $0 < Z < 1$ , representing the fact that some spectral weight is transferred to an incoherent background. Despite this renormalization, the defining geometric feature of the Fermi gas is preserved: the discontinuity in the occupation number. While the jump at the Fermi surface is reduced from 1 to  $Z$ , the Fermi surface itself remains sharp at  $T = 0$ . Furthermore, Luttinger's Theorem guarantees that the volume enclosed by the Fermi surface is invariant under interactions, provided the system remains a Fermi liquid. Thus, the interacting system retains the essential

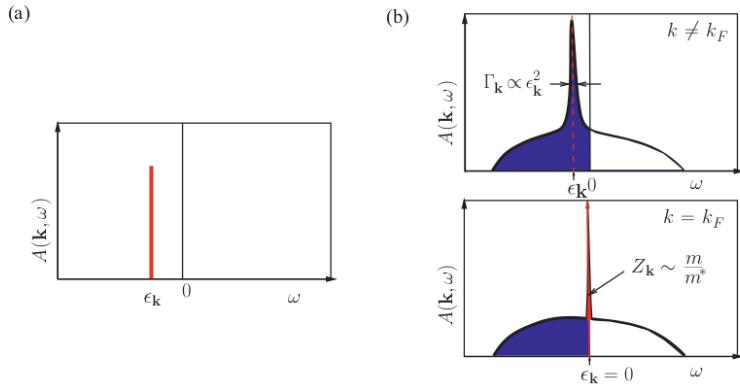
structure of a Fermi surface and well-defined quasiparticles, providing the foundation for Landau's phenomenological energy functional and its experimental predictions.

In a non-interacting Fermi gas at zero temperature, the particle momentum distribution is a sharp step function. All the states with energy below  $\mu$  are occupied, meaning all states above are empty. This sharp Fermi surface reflects the fact that adding a particle creates an exact eigenstate of the system with infinite lifetime  $\tau$ . Landau's claim is that this structure survives interactions in a precise sense. For a Fermi liquid, the low energy excitations are quasiparticles whose occupation numbers obey the same step-function structure as the free gas:

$$\langle \psi | (\hat{n}_{\mathbf{k},\sigma})_{qp} | \psi \rangle = \theta(\mu - \epsilon_k)$$

where  $(\hat{n}_{\mathbf{k},\sigma})_{qp} = a_{\mathbf{k}\sigma}^\dagger a_{\mathbf{k}\sigma}$  counts the quasiparticles rather than bare electrons. Since now we don't have a perfect step function anymore for the physical electron momentum distribution, we need to rely on the renormalization of the spectral weight to find out  $\langle \hat{n}_k \rangle$ . We see that the momentum distribution will exhibit a discontinuity of height  $Z$  at the Fermi surface, rather than a jump of height 1 as in the free Fermi gas.

Below is the figure illustrating this from Coleman. The spectral function consists of a sharp delta function at the single particle energy. In an interacting Fermi liquid, we have a narrow quasiparticle peak. For momenta away from the Fermi surface, this quasiparticle peak is broadened, reflecting a finite lifetime. But as  $k \rightarrow k_F$ , the width of this peak vanishes, meaning that the quasiparticle lifetime diverges and we get an infinite lifetime. The survival of this jump is one of the defining signs of a Fermi liquid. Despite strong interactions such as electron-electron interactions, the system still retains a sharply defined Fermi surface and well-defined quasiparticles.



(a) In a non-interacting Fermi system, the spectral function is a sharp delta function at  $\omega = \epsilon_k$ . (b) In an interacting Fermi liquid for  $k \neq k_F$ , the quasiparticle forms a broadened peak of width  $\Gamma_k$  at  $\omega_k$ . If  $k = k_F$ , this peak becomes infinitely sharp, corresponding to a long-lived quasiparticle on the Fermi surface. The weight in the quasiparticle peak is  $Z_k \sim m/m^*$ , where  $m^*$  is the effective mass.

## 2 Why Quasiparticle survive

We have asserted that the quasiparticle spectral peak remains sharp near the Fermi surface, implying an infinite lifetime at the Fermi energy. However, in an interacting system, we generally expect excited states to decay. A quasiparticle with energy  $\epsilon_k > \mu$

should scatter off the particles in the filled Fermi sea, losing energy and momentum. We must explain why this scattering rate  $\Gamma \rightarrow 0$  as  $\epsilon \rightarrow \mu$ .

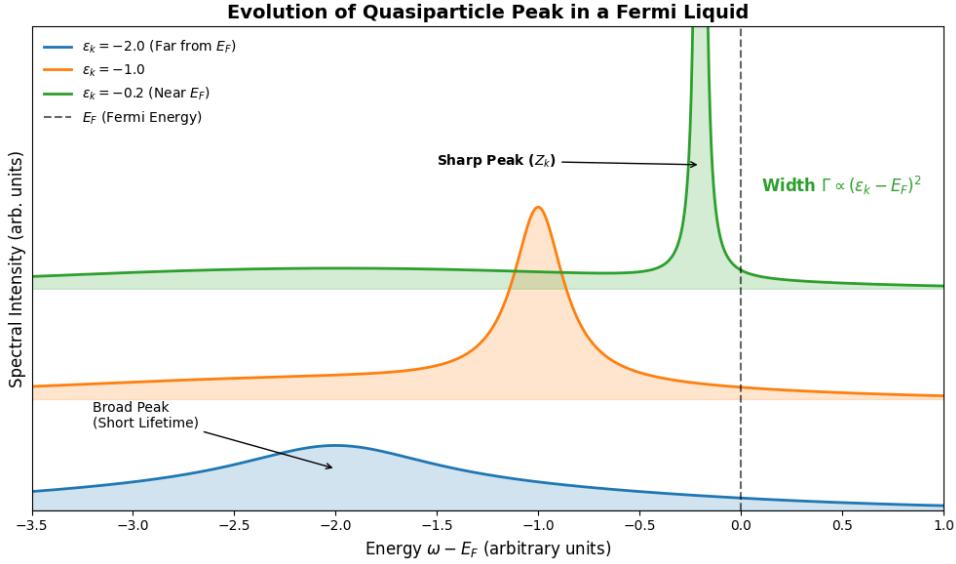


Figure 1: Evolution of the quasiparticle spectral peak in a Landau Fermi Liquid. Excitations far from the Fermi energy show a broad spectral peak corresponding to a short quasiparticle lifetime. As the excitation energy approaches the Fermi energy  $E_F$ , the peak sharpens and the width  $\Gamma$  decreases to 0. This reflects the divergence of the quasiparticle lifetime near the Fermi surface and underlies the stability of the Fermi liquid.

To formally show the stability of the Quasiparticle, one can follow [1] in section 6.7.2. In this paper, we detail the scattering mechanism and motivate the derivation heuristically and state the final result. Consider a quasiparticle in state  $|1\rangle$  with energy  $\epsilon_1 > 0$  (measured relative to the Fermi energy). The dominant decay channel is a three-body scattering event where the incoming quasiparticle interacts with a particle in the Fermi sea (state  $|2\rangle$ ,  $\epsilon_2 < 0$ ), scattering them into two new quasiparticle states  $|3\rangle$  and  $|4\rangle$  outside the Fermi sea. This process is denoted  $1 + 2 \rightarrow 3 + 4$ .

The amplitude to produce an outgoing hole in state  $\bar{2}$  is equal to the amplitude to absorb an incoming particle in state  $2$ , so denote:

$$a(1 \rightarrow \bar{2} + 3 + 4) = a(1 + 2 \rightarrow 3 + 4) \equiv a(1, 2; 3, 4).$$

The rate of this decay,  $\Gamma$ , is given by Fermi's Golden Rule. Summing over all possible collision partners and final states:

$$\Gamma \propto \sum_{2,3,4} |a(1, 2; 3, 4)|^2 \delta(\epsilon_1 + \epsilon_2 - \epsilon_3 - \epsilon_4) \delta(\mathbf{k}_1 + \mathbf{k}_2 - \mathbf{k}_3 - \mathbf{k}_4) \times (\text{Occupancy Factors})$$

where  $a$  is the effective interaction matrix element.

As Coleman describes it in [1], the crucial physics lies in the occupancy factors imposed by the Pauli Exclusion Principle. We have that state 2 must be initially occupied (i.e part of the Fermi sea):  $\epsilon_2 < 0$ . States 3 and 4 must be initially empty (i.e final states above the Fermi sea):  $\epsilon_3 > 0, \epsilon_4 > 0$ . Combining these with the energy conservation requirement  $\epsilon_1 + \epsilon_2 = \epsilon_3 + \epsilon_4$ , we find strict limits on the available energies. Since  $\epsilon_3, \epsilon_4 > 0$ , we must have:

$$\epsilon_3 + \epsilon_4 = \epsilon_1 - |\epsilon_2| < \epsilon_1$$

This implies that the excitation energy  $\epsilon_1$  must be shared between the two final quasiparticles. Consequently, all participating energies are constrained to a narrow shell of width  $\epsilon_1$  around the Fermi surface.

Following the detailed calculation of this three-body decay process (see Coleman), one finds the intermediate step of the energy dependence of the decay rate is governed by a phase-space integral of the form:

$$\Gamma \propto \int d\epsilon_2 \int d\epsilon_3 \int d\epsilon_4 \delta(\epsilon_1 + \epsilon_2 - \epsilon_3 - \epsilon_4)$$

and finally Coleman arrives at the result for the scattering of a spinless (spin only modifies the numerical prefactors) quasiparticle as:

$$\Gamma = \frac{1}{\tau} \propto (\epsilon^2 + (\pi k_B T)^2)$$

This result is one of the most vital results of the Fermi Liquid Theory. At the Fermi Surface ( $\epsilon \rightarrow \mu$ ), we have at zero temperature, the scattering rate vanishes quadratically. This means the lifetime  $\tau \sim 1/\epsilon^2$  diverges. The quasiparticle becomes an asymptotically exact eigenstate, justifying the sharp Fermi surface. For a finite temperature, at  $\epsilon = \mu$  but finite  $T$ , the lifetime behaves as  $\tau \propto T^{-2}$ . This finite lifetime at non-zero temperature is directly responsible for the  $T^2$  dependence of the electrical resistivity,  $\rho \propto T^2$ , a hallmark experimental signature of Fermi liquids.

### 3 The Landau Energy Functional

Having established the stability of quasiparticles, we now require a formalism to calculate the macroscopic properties of the Fermi liquid. Landau's central insight was that since quasiparticles are well-defined near the Fermi surface, the total energy of the system is a functional of their distribution function  $n_{\mathbf{p}\sigma}$ . We consider a small deviation  $\delta n_{\mathbf{p}\sigma}$  from the ground state distribution. Expanding the total energy density  $E$  to second order in  $\delta n$ , we obtain the fundamental equation of Landau Fermi Liquid theory:

$$E = E_0 + \sum_{\mathbf{p}\sigma} \epsilon_{\mathbf{p}}^{(0)} \delta n_{\mathbf{p}\sigma} + \frac{1}{2} \sum_{\mathbf{p}\sigma, \mathbf{p}'\sigma'} f_{\mathbf{p}\sigma, \mathbf{p}'\sigma'} \delta n_{\mathbf{p}\sigma} \delta n_{\mathbf{p}'\sigma'}$$

The first order term,  $\epsilon_{\mathbf{p}}^{(0)}$ , represents the energy of a single independent quasiparticle. The coefficient

$$\epsilon_{\mathbf{p}}^{(0)} \equiv E_{\mathbf{p}\sigma}^{(0)} - \mu = \frac{\delta E}{\delta n_{\mathbf{p}\sigma}}$$

describes the excitation energy of an isolated quasiparticle. Given that we can ignore the spin-orbit interactions, then we find the total magnetic moment is a conserved quantity, so the magnetic moments of the quasiparticles are preserved by interactions. If we expand the quasiparticle energy linearly in momentum near the Fermi surface:

$$E_p^{(0)} = v_F(p - p_F) + \mu^{(0)}$$

with  $v_F$  being Fermi velocity at the Fermi energy  $\mu^{(0)}$ , with  $\mu^{(0)}$  being the chemical potential in the ground state. We can then define the quasiparticle effective mass  $m^*$  in terms of  $v_F$ :

$$v_F = \left. \frac{d\epsilon_p^{(0)}}{dp} \right|_{p=p_F} = \frac{p_F}{m^*}$$

This result can be used to define the density of states for the quasiparticle at  $\epsilon = 0$ :

$$\mathcal{N}^* = \frac{m^* p_F}{\pi^2 \hbar^3}$$

This way, the effective mass determines the density of states at the Fermi energy. Large effective masses lead to large densities of states.

The second term introduces the Landau Interaction Function,

$$f_{\mathbf{p}\sigma,\mathbf{p}'\sigma'} = \left. \frac{\delta^2 E}{\delta n_{\mathbf{p}\sigma} \delta n_{\mathbf{p}'\sigma'}} \right|_{\delta n_{\mathbf{p}''\sigma''}=0}$$

describes the interactions between the quasiparticle at the Fermi surface. The partials are calculated in a "Frozen" Fermi sea, which is akin to treating the other quasiparticle occupancies fixed.

### 3.1 The Landau Parameters

Due to the ability to parametrize the interactions in terms of a small number of parameters called the Landau Parameters, we see the true power of the Landau Fermi liquid theory. We are able to use these parameters to renormalize the non-interacting Fermi liquid theory by the feedback effect of interactions on quasiparticle energies. In a Landau Fermi liquid where spin is conserved, we see that the interaction is invariant under spin rotations and can be written in general as:

$$f_{\mathbf{p}\sigma,\mathbf{p}'\sigma'} = f_{\mathbf{p}\mathbf{p}'}^s + f_{\mathbf{p}\mathbf{p}'}^a \sigma \sigma'$$

where the spin-dependent term of the interaction is the magnetic component of the quasiparticle interaction.

For an isotropic Fermi liquid (such as liquid  ${}^3\text{He}$  or a simple metal), the interaction function depends only on the relative angle between momenta and the relative spin orientation. We can thus decompose  $f_{\mathbf{p}\mathbf{p}'}$  into spin-symmetric ( $f^s$ ) and spin-antisymmetric ( $f^a$ ) components. To see this, we first write:

$$f_{\mathbf{p}\mathbf{p}'}^{s,a} = f^{s,a}(\cos \theta)$$

Then convert the interaction to a dimensionless function by multiplying by density of states  $\mathcal{N}$ .

$$F^{s,a}(\cos \theta) = \mathcal{N} f^{s,a}(\cos \theta)$$

Then we expand the functions as a multipole expansion in terms of Legendre polynomials:

$$F^{s,a}(\cos \theta) = \sum_{l=0}^{\infty} (2l+1) F_l^{s,a} P_l(\cos \theta)$$

The coefficients here,  $F_l^s, F_l^a$  are the **Landau Parameters**. The spin-symmetric components  $F_l^s$  account for the non-magnetic part of the interaction, while the spin-antisymmetric  $F_l^a$  define the magnetic component of the interaction [1]. Given in Coleman eqn 6.41, we find a convenient form to calculate these parameters:

$$F_l^{s,a} = 2 \sum_{\mathbf{p}'} f_{\mathbf{p},\mathbf{p}'}^{s,a} P_l(\cos \theta_{\mathbf{p},\mathbf{p}'}) \delta(\epsilon_{\mathbf{p}'})$$

## 4 Predictions & Observables

Using this model, we can derive exact expressions for experimentally measurable quantities. The strong interactions do not change the functional form of the response (e.g., linear specific heat), but they do rescale the coefficients [1].

Firstly, we have the **Effective Mass** ( $m^*$ ). The interaction between a moving quasiparticle and the surrounding fluid creates a backflow that alters its inertia. This is encoded by the  $F_1^s$  parameter:

$$\frac{m^*}{m} = 1 + \frac{F_1^s}{3}$$

In strongly correlated systems like Heavy Fermions,  $F_1^s$  can be large and positive, leading to quasiparticles that are hundreds of times heavier than bare electrons.

We also have the **Electronic Specific Heat** ( $C_V$ ). We know that low-temperature specific heat of a metal is linear,  $C_V = \gamma T$ . With  $\gamma = \frac{\pi^2}{3} k_B^2 \mathcal{N}$ . Since  $\gamma$  is proportional to the density of states, which in turn is linear in mass ( $N(0) \propto m^*$ ), the specific heat is enhanced by the same factor:

$$C_V = \frac{m^*}{m} C_V^{free} = \gamma_{free} \left( 1 + \frac{F_1^s}{3} \right) T$$

This explains why strongly interacting metals still exhibit the linear specific heat characteristic of free electrons, but with a significantly larger magnitude.

The Pauli paramagnetic **Spin Susceptibility** ( $\chi_s$ ) measures the response to a magnetic field. It is renormalized by the spin-antisymmetric parameter  $F_0^a$ :

$$\chi_s = \frac{\mu_F^2 \mathcal{N}}{1 + F_0^a}$$

This result is vital for understanding magnetic instabilities. If interactions are strongly ferromagnetic ( $F_0^a \rightarrow -1$ ), the susceptibility diverges, indicating a phase transition to a ferromagnetic state (the Stoner instability).

The charge susceptibility is given by:

$$\chi_c = \frac{\mathcal{N}}{1 + F_0^s}$$

and is related to the **Compressibility/Bulk Modulus** ( $\kappa$ ) which measures the resistance of the liquid to density changes. For a neutral fluid the bulk modulus is related to the charge susceptibility per unit volume as:

$$\kappa = \frac{n^2}{\chi_c} \quad n = \frac{N}{V}$$

with  $n$  is particle density. And so a smaller charge susceptibility implies a stiffer, less compressible fluid than the free gas.

## 5 Conclusion

To conclude, we end the paper saying that the predictions of the Landau Fermi Liquid theory models real experiments in a wide range of physical systems. In simple metals

such as copper [2], the electronic specific heat is observed to be linear in temperature at low  $T$ , with a coefficient  $\gamma$  consistent with a renormalized effective mass.

One of the clearest realization of a Fermi liquid is a liquid  $^3\text{He}$  where precise measurements of the specific heat, spin susceptibility and compressibility agree quantitatively with Landau's theory. More recently, angle-resolved photoemission spectroscopy (ARPES) experiments in simple metals and correlated electron systems have directly observed long-lived quasiparticle peaks near the Fermi Surface, providing direct evidence of the quasiparticle concept underlying the Fermi Liquid theory. Such an example is the 'Quasi-particle spectra and Fermi surface mapping' by R. Classen. [3]

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