Fermi liquid theory: A brief survey in memory of Gerald E. Brown

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Abstract
I present a brief review of Fermi liquid theory, and discuss recent work on Fermi liquid theory in dilute neutron matter and cold atomic gases. I argue that renewed interest in transport properties of quantum fluids provides fresh support for Landau’s approach to Fermi liquid theory, which is based on kinetic theory rather than effective field theory and the renormalization group. I also discuss work on non-Fermi liquids, in particular dense quark matter.

1. Introduction
One of Gerry’s main scientific pursuits was to understand the nuclear few and many-body problem in terms of microscopic theories based on the measured two and three-nucleon forces. One of the challenges of this program is to understand how the observed single-particle aspects of finite nuclei, in particular shell structure and the presence of excited levels which carry the quantum numbers of single particle states, can be reconciled with the strong nucleon-nucleon force, and how single particle states can coexist with collective modes. A natural framework for addressing these questions is the Landau theory of Fermi liquids. Landau Fermi liquid theory describes a possibly strongly correlated, Fermi system which is adiabatically connected to a free Fermi gas. In particular, the system has a Fermi surface, and the excitations are quasi-particles with the quantum numbers of free fermions, but with modified dispersion relations and effective interactions. These quasi-particles coexist with collective modes, for example zero sound.

Gerry reviewed Fermi liquid theory in a number of his books and other writings. The conference that celebrated his 60th birthday was titled “Windsurfing the Fermi Sea” [1]. In the introduction of Unified Theory of Nuclear Models and Forces (3rd edition, 1970) Gerry writes:

Many improvements could have been made, especially in Chapter XIII on effective forces in nuclei, but time is short, and I shall make them in later editions, when I am too old to ski. Of course, nobody will be interested in the subject by then.

This prediction turned out to be incorrect. In his final decade at Stony Brook Gerry trained and mentored a remarkable group of students who have helped to reinvigorate the study of effective forces in nuclei [2,3].
2. Landau Fermi liquid theory

Consider a cold Fermi system in which the low energy excitations are spin 1/2 quasi-particles. Landau proposed to define a distribution function $f_p = f^0_p + \delta f_p$ for the quasi-particles. Here, $f^0_p$ is the ground state distribution function, and $\delta f_p \ll f^0_p$ is a small correction. The energy density can be written as \[ E = E_0 + \int d\Gamma_p \frac{\delta E}{\delta f_p} \delta f_p + \frac{1}{2} \int d\Gamma_p d\Gamma_p' \frac{\delta^2 E}{\delta f_p \delta f_p'} \delta f_p \delta f_p' + \ldots , \] (1) with $d\Gamma_p = d^3 p/(2\pi)^3$. Functional derivatives of $E$ with respect to $f_p$ define the quasi-particle energy $E_p$ and the effective interaction $t_{pp'}$

Near the Fermi surface we can write $E_p = v_F (|\vec{p}| - \pi_F)$, where $v_F$ is the Fermi velocity, $p_F$ is the Fermi momentum, and $m^* = p_F/v_F$ is the effective mass. We can decompose $t_{pp'} = F_{pp'} + G_{pp'} \vec{\sigma}_1 \cdot \vec{\sigma}_2$. On the Fermi surface the effective interaction is only a function of the scattering angle and we can expand the angular dependence as \[ F_{pp'} = \sum_l F_l (\cos \theta_{\vec{p} \vec{p}'}), \] (3) where $P_l(x)$ is a Legendre polynomial, and $G_{pp'}$ can be expanded in an analogous fashion. The coefficients $F_l$ and $G_l$ are termed Landau parameters.

The distribution function satisfies a Boltzmann equation \[ \left( \partial_t + \vec{v}_p \cdot \vec{\nabla} + \vec{F}_p \cdot \vec{\nabla} \right) f_p(x, t) = C[f_p] \] (4) where $\vec{v}_p = \vec{v}_p E_p$ is the quasi-particle velocity, $\vec{F}_p = -\vec{\nabla}_x E_p$ is an effective force, and $C[f_p]$ is the collision term. Conserved currents can be defined in terms of $f_p$ and the single particle properties $E_p$ and $v_p$. For example, we can write the mass density $\rho$ and mass current $\vec{j}$ as \[ \rho = \int d\Gamma_p m f_p, \quad \vec{j} = \int d\Gamma_p m \vec{v}_p f_p, \] (5) where $d\Gamma_p = d^3 p/(2\pi)^3$. The Boltzmann equation implies that the current is conserved, $\partial_t \rho + \vec{\nabla} \cdot \vec{j} = 0$. The conditions given in equ. (2) play an important role in proving conservation laws for energy and momentum, and in establishing sum rules.

A different approach to Fermi liquid theory was popularized by Polchinski [5] and Shankar [6], see also [7,8]. Consider free non-relativistic quasi-particles near Fermi surface. The system is described by the action \[ S_{FL} = \int dt \int \frac{d^3 p}{(2\pi)^3} \psi(p) \left( i\partial_t - \left[ \epsilon(p) - \epsilon_F \right] \right) \psi(p) \] (6) Near the Fermi surface we can expand the momentum $\vec{p} = \vec{p}_F + \vec{t}$ and \[ \epsilon(p) - \epsilon_F = \vec{v}_F(k) \cdot \vec{t} + O(\vec{t}^2) \] (7)
We are interested in the question whether \( S_{FL} \) is a possible fixed point of the renormalization group. For this purpose we study the scaling behavior of the action as \( l \to s l \). The free fermion action is invariant under this rescaling provided we assign the following scaling dimensions

\[
[k] = 0, \quad [l] = 1, \quad [\partial_l] = 1, \quad [d^3 p] = 1, \quad [\psi] = -\frac{1}{2}.
\]  

Consider now the effect of a four-fermion interaction

\[
S_{int} = \int dt \left[ \prod_{i=1}^4 \int \frac{d^3 p_i}{(2\pi)^3} \right] \psi^\dagger(p_3)\psi^\dagger(p_2)\psi(p_1)\delta^3(p_{tot})U(p_1, p_2, p_3, p_4) .
\]

For generic values of the momenta we can use the scaling dimensions in equ. (8) to establish that interaction terms are irrelevant near the Fermi surface. An exception can occur if the dominant components of the momenta cancel and the delta function constrains the small components \( l_i \). There are two configurations for which this occurs. One is BCS scattering \((\hat{p}_F, -\hat{p}_F) \to (\hat{p}_F, -\hat{p}_F)\) which is characterized by the interaction

\[
U(-\hat{p}_3, \hat{p}_3, -\hat{p}_1, \hat{p}_1) = \sum_l V_l \hat{p}_l \cdot \hat{p}_3 ,
\]

where \( P_l(x) \) are Legendre polynomials. The \( V_l \) are marginal interactions. If any \( V_l \) is attractive then loop corrections will grow logarithmically as \( s \to 0 \), and lead to BCS superfluidity. The other configuration is generalized forward scattering, where \( \hat{p}_1 \cdot \hat{p}_2 = \hat{p}_3 \cdot \hat{p}_4 \), which is characterized by

\[
U(\hat{p}_4, \hat{p}_3, \hat{p}_2, \hat{p}_1) = F(\hat{p}_1 \cdot \hat{p}_2, \phi_{12, 34}) + G(\hat{p}_1 \cdot \hat{p}_2, \phi_{12, 34}) \sigma_1 \cdot \sigma_2 .
\]

We can write

\[
F(z, 0) = \sum_l F_l P_l(z) ,
\]

and \( F_l \) are the Landau Fermi liquid parameters defined above. The Fermi liquid parameters remain marginal even if loops are included.

### 3. Neutron matter

In the case of neutron matter the renormalization group evolution towards the Fermi surface was carried out explicitly by Schwenk, Friman, and Brown [9], see Fig. 1. Schwenk et al. use \( V_{sym} \) as the bare interaction at the UV scale \( \Lambda = \sqrt{2} p_F \). In nuclear matter three-nucleon forces are important, but there is no local \( s \)-wave interaction between three neutrons, and as a consequence three nucleon forces are much less important in neutron matter. Schwenk et al. employ the functional renormalization group equation in order to evolve the interaction to the Fermi surface. The resulting Fermi liquid parameters \( F_0 \) and \( F_1 \) are shown in Fig. 1. The parameter \( F_1 \) is related to the effective mass \(^4\)

\[
\frac{m^*}{m} = 1 + \frac{F_1}{3} .
\]

\(^4\)This is a curious relation from the point of view of effective field theory, as it relates parameters of the EFT to a bare parameter. First of all, we note that both \( m^* \) and \( F_1 \) are coefficients of marginal operators, so they can indeed occur together. Second, we can view this relation as a non-renormalization theorem that follows from Galilean invariance. Indeed, the bare mass is part of the Galilean algebra.
We observe that in the regime $0.6 \text{ fm}^{-1} < p_F < 1.4 \text{ fm}^{-1}$ the effective mass exceeds unity. It is interesting to compare this result to recent data from cold atomic Fermi gases at unitarity. The unitary Fermi gas realizes the limit $a \to \infty$, where $a$ is the $s$-wave scattering length between the fermions. This system is an interesting model system for dilute neutron matter, because the $nn$ scattering length is anomalously large, $a_{nn} \approx -19 \text{ fm}$. Analyzing thermodynamic observables of the unitary Fermi gas just above the superfluid transition, Nascimbene et al. find $m'/m = 1.13 \pm 0.03$ [12], consistent with the sign of $F_1$ in Fig. 1, but somewhat larger in magnitude.

4. Non-Fermi liquid effective field theory

Effective field theory ideas can also be applied to dense quark matter. As we will see, quark matter is not a Fermi liquid in the strict sense, but many of the ideas of Fermi liquid theory survive. In particular, quark matter above the critical temperature for color superconductivity has quasi-particle and quasi-hole excitations that carry the quantum numbers of quarks.

QCD is a gauge theory, and the main new ingredient in an effective field theory of dense quark matter is the coupling of quarks to transverse gauge fields. This coupling is of the form $\vec{v} \cdot \vec{A}$, where $\vec{v}$ is the velocity of the quark. At high density the Fermi momentum is large, and the emission of a low momentum gluon cannot change the velocity of the quark. An effective theory of quasi-quarks and quasi-holes interacting with soft gluons can be constructed by covering the Fermi surface with patches labeled by the local Fermi velocity $\vec{v}$. The effective lagrangian is given by [13, 14, 15].

\begin{equation}
\mathcal{L} = \sum_v \psi_v^\dagger \left( i v \cdot \partial - \frac{1}{2p_F} D_v^2 \right) \psi_v + \mathcal{L}_{4f} + \mathcal{L}_{HDL} = \frac{1}{4} G_{\mu\nu} G^{\mu\nu} + \ldots ,
\end{equation}

This idea is discussed in [6] as the basis of a "large N" approach to ordinary Fermi liquid theory.
Figure 2: One-loop contributions to the quark self energy and the quark-gluon vertex. Near the Fermi surface the Feynman diagrams scale as $\omega \log(\omega)$, $\omega^{1/3}$ and $\omega^{2/3}$, respectively.

where $v_\mu = (1, \vec{v})$ if the four-velocity. The field $\psi_v$ describes particles and holes with momenta $p = \mu (0, \vec{v}) + k$, where $k \ll \mu$. We can decompose $k = k_0 + k_\parallel + k_\perp$ with $\vec{k}_\parallel = \vec{v}(\vec{k} \cdot \vec{v})$ and $\vec{k}_\perp = \vec{k} - \vec{k}_\parallel$.

Hard gluon exchanges are described by a local four-fermion interaction $L_{4f}$. This interaction has the same structure as ordinary Fermi-liquid theory. Hard fermion loops are absorbed into the hard dense loop lagrangian $L_{HDL}$. This interaction is somewhat subtle, because it cannot be written in terms of a local lagrangian. Braaten and Pisarski showed that

$$L_{HDL} = -\frac{m^2}{2} \int \frac{d\vec{v}}{4\pi} G^\mu_\alpha \frac{\vec{v}^\nu v^\theta}{(v \cdot D)^2} G^\nu_\beta,$$

where $m^2 = N_f g^2 \mu^2 / (4\pi^2)$ is the dynamical gluon mass and $\hat{v}$ is a unit vector. The hard dense loop action describes static screening of electric gauge fields and dynamic screening of magnetic modes. Since there is no static magnetic screening we find that low energy gluon exchange is dominated by magnetic modes. The transverse gauge boson propagator is given by

$$D_{ij}(k) = \frac{\delta_{ij} - \hat{k}_i \hat{k}_j}{k_0^2 - \hat{k}^2 + i\frac{\pi}{2}m^2 \frac{k_0}{|\hat{k}|}},$$

where we have assumed that $|k_0| < |\hat{k}|$.

The gluon propagator is dominated by modes with momenta $|\hat{k}| \sim (m^2 k_0)^{1/3} \gg k_0$. This relation leads to anomalous scaling relations as we approach the Fermi surface. Consider a generic Feynman diagram and scale all energies by a factor $s$. From the relation quoted above we conclude that gluon momenta scale as $|\vec{k}| \sim s^{1/3}$, which implies that the momentum of a gluon is parametrically large to its energy, and that gluons are far off-shell and space-like.

The quark dispersion relation is $k_0 \sim s$, $k_\parallel \sim s^{2/3}$, $k_\perp \sim s^{1/3}$.

and $k_0 \ll k_\parallel \ll k_\perp$. In this regime the quark and gluon propagators are given by

$$S_{\alpha\beta}(p) = \frac{i\delta_{\alpha\beta}}{p_0 - p_\parallel - \frac{p_\perp^2}{2} + i\epsilon \text{sgn}(p_0)}; \quad D_{ij}(k) = \frac{-i\delta_{ij}}{k_0^2 - k_\perp^2 - \frac{m^2 k_0}{|\hat{k}|}},$$

where $\vec{k} = \vec{v}(\vec{k} \cdot \vec{v})$ and $\vec{k} = \vec{k} - \vec{k}_\parallel$. The quark dispersion relation is

$$k_0 \sim s, \quad k_\parallel \sim s^{2/3}, \quad k_\perp \sim s^{1/3},$$

and $k_0 \ll k_\parallel \ll k_\perp$. In this regime the quark and gluon propagators are given by
Figure 3: This figure shows the one-loop correction to the quark-quark interaction in the BCS channel (left panel) and vertex of an external current (right panel). Both diagrams are kinematically enhanced, and scale as \( \log^2(\omega) \) and \( \log(\omega) \), respectively.

and the quark gluon vertex is \( gv(\lambda^a/2) \). Higher order terms can be found by expanding the quark and gluon propagators as well as the HDL vertices in powers of the small parameter \( \epsilon \equiv \omega/m \).

We can compare the scaling behavior of the interaction to the result for a Fermi liquid, see equ. (8). Consider the scale transformation \((x_0, x_{||}, x_\perp) \rightarrow (s^{-1}x_0, s^{-2/3}x_{||}, s^{-1/3}x_\perp)\). The fields scale as \( \psi \rightarrow s^{5/6}\psi \) and \( A_i \rightarrow s^{5/6}A_i \). We find that the scaling dimension of all interaction terms is positive. The quark gluon vertex scales as \( s^{1/6} \), the HDL three gluon vertex scales as \( s^{1/2} \), and the four gluon vertex scales as \( s \). Since higher order diagrams involve at least one pair of quark gluon vertices the expansion involves positive powers of \( s^{1/3} \). The low energy theory is indeed weakly coupled, but the expansion involves fractional powers (and logarithms) of the energy.

The effective field theory has been used to study a wide range of problems:

1. The quark self energy scales as \( \Sigma(\omega) \sim \log(\omega) \), see Fig. 2. This result implies that the Fermi velocity and the wave function renormalization go to zero as \( \omega \rightarrow 0 \). The logarithm in \( \Sigma(\omega) \) also leads to logarithmic terms in the specific heat \([17]\). All of these effects are in contrast to a Fermi liquid. In perturbation theory \( \Sigma(\omega) \sim g^2 \log(\omega) \), and there are claims in the literature that perturbation theory breaks down at the scale \( \omega \sim \exp(-1/g^2) \). This is not the case, as shown in \([18]\). We also find a QCD version of Luttinger’s theorem \([19]\). The quark density is given by the volume of the “Fermi surface”, defined by the condition that the inverse quark propagator \( S^{-1}(\omega=0, p) \) changes sign \([20]\).

2. Corrections to the quark-gluon vertex are dominated by the abelian diagram in Fig. 2. The scaling rules imply that

\[
\Gamma_\mu = gv_\mu(\lambda^a/2) \left( 1 + O(\epsilon^{1/3}) \right). \tag{19}
\]

This is a QCD version of Migdal’s theorem, which states that the renormalization of the electron-phonon vertex is suppressed by the ratio \( \sqrt{m/M} \), where \( m \) is the mass of the electron and \( M \) is the mass of the ions \([21]\). This factor \( \sqrt{m/M} \) is analogous to the small parameter in dense QCD, because it governs the ratio of the phonon velocity to the Fermi velocity of the electrons, which determines the ratio of the typical electron and phonon momenta. As a consequence of Migdal’s theorem we can compute the superfluid gap using an approximation which includes loop correction to the gluon propagator, but does include vertex corrections.

3. In ordinary Fermi liquid theory, there are certain kinematical regimes, the BCS and zero sound channels, in which the interaction is enhanced, see equ. (10,11). The same situation
arises in non-Fermi liquid quark matter. Consider rescattering in the BCS channel as shown in Fig. 3. The diagram scales as $g^2 \log^2(\omega)$, and the interaction becomes non-perturbative at the scale $\omega \sim \exp(-1/g)$. This effect leads to color superconductivity, and a parametrically large gap $\Delta \sim \mu \exp(-1/g)$ [22]. A similar phenomenon occurs when we consider the vertex of an external gauge field with coupling $e$. The one-loop correction in the regime of small time-like momenta is $\Gamma_\mu \sim eg^2 v_\mu \log(\omega)$, and non-perturbative effects become important for $\omega \sim \exp(-1/g^2)$ [23].

5. Transport theory

The examples discussed in the previous two sections are based on the “modern”, EFT based, approach to Fermi liquid theory. In this section I would like to argue that recent calculations of transport properties of strongly correlated Fermi liquids have reinvigorated Landau’s original approach, which is based on kinetic theory. As an illustration I will discuss recent work on the bulk viscosity of a dilute Fermi gas near unitarity [24]. Interactions between quasi-particles play an essential role in this problem, because one can show that the bulk viscosity is zero if the “interaction measure”

$$\Lambda = E - \frac{3}{2} P$$

vanishes. In kinetic theory we view $E$ as a functional of the distribution function $f_p$. The momentum density and the stress tensor are given by

$$\tilde{\pi}(\vec{x}, t) = \int d\Gamma_p \tilde{p}_p f_p(\vec{x}, t),$$

$$\Pi^{ij}(\vec{x}, t) = \int d\Gamma_p p^i p^j f_p(\vec{x}, t) + \delta^{ij} \left( \int d\Gamma_p E_p f_p(\vec{x}, t) - E(\vec{x}, t) \right).$$

Consistency with the laws of fluid dynamics requires that the momentum density is equal to the mass current,

$$\tilde{\pi}(\vec{x}, t) = \int d\Gamma_p m v^i p^j f_p(\vec{x}, t),$$

and that the momentum density is conserved,

$$\partial_0 \tilde{\pi}(\vec{x}, t) + \nabla j^{ij}(\vec{x}, t) = 0.$$ \hspace{1cm} (24)

These relation can only be satisfied if $E_p = (\delta E)/(\delta f_p)$, which is the fundamental relation in Landau’s approach to Fermi liquid theory.

In our work [24] we compute the equation of state and the quasi-particle properties as an expansion in the fugacity $z = \exp(\mu/T)$. This calculation is, strictly speaking, not a Landau Fermi liquid calculation because the fugacity expansion is reliable at high temperature, not at low temperature. However, since we employ a functional expansion of $E$, the calculation can be generalized to the Fermi liquid case by replacing the bare interaction with an effective interaction.

\footnote{At unitarity the fluid is exactly scale invariant, and the bulk viscosity vanishes.}
and by including the effects of quantum statistics in the collision term. Using the bare interaction the quasi-particle energy near unitarity is $E_p = \epsilon_p + \text{Re} \Sigma(p) + i \text{Im} \Sigma(p)$ with $\epsilon_p = p^2/(2m)$ and

$$
\text{Re} \Sigma(p) = -4 \sqrt{2}\zeta T \frac{1}{a \sqrt{mT}} \sqrt{\frac{T}{\epsilon_p}} F_D \left( \sqrt{\frac{\epsilon_p}{T}} \right), \tag{25}
$$

$$
\text{Im} \Sigma(p) = -2\zeta T \frac{1}{\sqrt{\epsilon_p}} E_r \left( \sqrt{\frac{\epsilon_p}{T}} \right), \tag{26}
$$

where $F_D$ is Dawson’s Integral, and $E_r$ is the error function. Because $\text{Im} \Sigma \sim \zeta T \ll \epsilon_p \sim T$ quasi-particles are well defined.

In an interacting fluid the total energy is the sum of a non-interacting term and the interaction energy. Bulk viscosity arises from the fact that in an expanding system the relative magnitude of these two contributions may deviate from the equilibrium value. We find

$$
\zeta = \frac{1}{96\pi^{5/2}} (mT)^{3/2} \left( \frac{\zeta \lambda}{a} \right)^2, \tag{27}
$$

where $\lambda$ is the thermal de Broglie wave length. This result is consistent with the estimate $\zeta \sim \eta \Delta^2$, where $\Delta$ is the interaction measure, and $\eta \sim (mT)^{3/2}$ is the shear viscosity.

6. Frontiers

Fermi liquid theory continues to be a very active area of many-body physics, and there are many new ideas that I cannot do justice to in a short article. One interesting area of research is to understand the way anomalies in the underlying quantum field theory are implemented in the Fermi liquid theory. Son and Yamamoto showed that anomalies give rise to non-zero Berry curvature of the Fermi surface \[25\]. The kinetic equation then leads to an anomalous conservation law,

$$
\partial_0 n + \vec{\nabla} \cdot \vec{j} = \pm \frac{1}{4\pi^2} \vec{E} \cdot \vec{B}, \tag{28}
$$

where the $\pm$ sign refers to left/right handed fermions, and $\vec{E}$, $\vec{B}$ are electric and magnetic $U(1)$ fields. Son and Yamamoto also showed how this equation arises in the high density effective field theory described in Sect.\[4\][26].

Another interesting area of research is the study of holographic Fermi and non-Fermi liquids \[27, 28\]. These constructions are based on charged black holes embedded in asymptotically AdS (Anti de Sitter) spaces. According to the AdS/CFT correspondence quantum gravity on this background is holographically dual to certain field theories in flat space. By adding spinor fields to the gravitational theory it is possible to realize holographic Fermi surfaces with a variety of Fermi liquid and non-Fermi liquid spectral functions. Faulkner et al. describe a holographic model that gives a retarded spinor propagators of the form \[28\].

$$
G_R(\omega, p) = \frac{h_1}{\omega - \nu_F(p - p_F) - \Sigma(\omega, p)}, \quad \Sigma(\omega, p) = c\omega^{2\nu_F}, \tag{29}
$$

where $h_1, \nu_F, c$ are constants. The model can realize $\nu_F = 1/2$, corresponding to a marginal Fermi liquid such as the one described in Sect.\[4\] as well as $\nu_F > 1/2$ and $\nu_F < 1/2$. An
important aspect of the AdS/CFT correspondence is that it is fairly straightforward to compute transport properties at strong coupling. As a result, one can relate quasi-particle properties, in particular the behavior of the fermion self energy near the Fermi surface, to transport properties, such as the conductivity spectral function.

7. Final Remarks

I met Gerry in 1991 at a meeting in Peniscola, Spain. At the meeting, Gerry recruited me for long runs along the beach. I came to Stony Brook as a postdoc in 1992 and stayed until 1995. I later returned as an Assistant Professor from 2000-2002. In the end, I only wrote one paper with Gerry [29], but the many lunch discussions, and the way Gerry attracted and mentored his students and postdocs made a lasting impression on me. I will miss the unusual combination of his dry and sometimes harsh wit with the deep concern he showed for the welfare of those close to him. Gerry liked to talk about his “eagles”[30]. Breit, Bethe, and Peierls. Just like these great scientists, Gerry appears to belong to a different era, one in which individual scientists could touch more people, and leave a more lasting impact than seems possible today.

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