The Quantum-Classical Metal

David G. Clarke, S. P. Strong, P. M. Chaikin, E. I. Chashechkina

In a normal Fermi liquid, Landau's theory precludes the loss of single-fermion quantum coherence in the low-energy, low-temperature limit. For highly anisotropic, strongly correlated metals, there is no proof that this remains the case, and quantum coherence for transport in some directions may be lost intrinsically. This loss of coherence should stabilize an unusual, qualitatively anisotropic non-Fermi liquid, separated by a zero-temperature quantum phase transition from the Fermi liquid state and categorized by the unobservability of certain interference effects. There is compelling experimental evidence for this transition as a function of magnetic field in the metallic phase of the organic conductor $(TMTSF)_2PF_6$ (where TMTSF is tetramethyltetraselenafulvalene).

At the heart of the modern understanding of metals is Landau's Fermi liquid theory (FLT) (*1*). According to this theory, the ground state and low-lying excited states of a system of interacting fermions can be placed into one-to-one correspondence with similar states in an appropriately chosen model of noninteracting fermions (such as a free-electron gas). For the interacting system, these low-lying states are labeled by occupation numbers for fermionic "quasiparticles," just as the states of the noninteracting system are labeled with the occupation numbers for the free fermions. In the quasiparticle basis, the ground state of the interacting system is a filled Fermi sea (*2*), and individual quasiparticles excited above this Fermi sea have a scattering rate from the other quasiparticles that vanishes much faster (quadratically) than their excitation energy (*3*). In this sense, the quasiparticles are weakly interacting, even though the interactions among the original fermions may have been quite strong.

This aspect of the theory accounts for (i) the otherwise mysterious successes of nearly free electron models of metals and, with the exception of various broken symmetry states that themselves at least fit into the Fermi liquid paradigm, (ii) why there were no serious experimental challenges to FLT until comparatively recently. In recent times, the cuprate superconductors (*4*) and other highly anisotropic metals (*5*) have exhibited properties very different from the predictions of FLT. We discuss a proposal for a breakdown of the theory for these highly anisotropic systems with strong electron-electron interactions.

One hallmark of a Fermi liquid is that quantum coherence for the motion of a single fermion in all directions is intrinsic; that is, it is present in the low-energy, low-temperature limit for pure systems. Here we present theoretical and experimental considerations strongly challenging the inevitability of such coherence for anisotropic metals with strong electron-electron interactions. In fact, the evidence for the loss of intrinsic quantum coherence for transport in one direction in the organic conductor $(TMTST)_2PF_6$ is compelling.

D. G. Clarke, P. M. Chaikin, and E. I. Chashechkina are with the Joseph Henry Laboratories of Physics, Princeton University, Post Office Box 708, Princeton, NJ 08544, USA. S. P. Strong is with the Institute for Advanced Study, Olden Lane, Princeton, NJ 08540, USA.

The Stability of Fermi Liquid Theory

For short-range interactions in two or more dimensions, FLT is valid at all orders in perturbation theory of the electron-electron interaction. This agreement results from the internal consistency of FLT: the filled Fermi sea of quasiparticles restricts the scattering phase space so thoroughly that excited quasiparticles experience arbitrarily small scattering rates for small enough excitation energies. As a result, the quasiparticles effectively become eigenexcitations for sufficiently small excitation energies, and their filled Fermi sea is an appropriate ground state. To find behaviors other than those expected in FLT, one must start outside of this self-consistent loop. One way of doing this is to begin with the exact solution of a system of interacting fermions in one dimension, where FLT generally fails.

Interacting fermions generally do not form Fermi liquids in one dimension because the phase space argument leading to the quadratically vanishing scattering rate is not valid in one dimension. Instead, the fermions in such models generically form Luttinger liquids. These liquids still have excitations with vanishing scattering rates, but the excitations are bosonic, rather than fermionic, in character and are related to the density wave modes of the underlying fermions (*6*). As a consequence, Luttinger liquids differ from Fermi liquids in a variety of ways. For example, the power laws occurring in many correlation functions are functions of the interaction strength. Moreover, the velocities of spin and charge excitations are generically different in Luttinger liquids, although they are always the same in Fermi liquids.

It is now well accepted that FLT fails for generic one-dimensional models and that Luttinger liquid behavior results. However, real materials are never one or even two dimensional; they are at most anisotropic, and therefore, a proposal for non-Fermi liquid behavior in a bulk material cannot be justified with one-dimensional results. Even a very "one-dimensional" material must be represented with coupled one-dimensional models. This coupling is, in principle, capable of restoring FLT at low enough energies. To argue that FLT fails in bulk metals, one must demonstrate that for some reason this coupling does not restore FLT. Such a demonstration requires the inclusion of, for example, single particle hopping between one-dimensional chains and a demonstration of something exotic about the character of the single particle hopping. In previous work (*7*) we have found evidence for a nontrivial breakdown of FLT related to a loss of what we termed the "coherence of interliquid hopping."

Quantum Coherence

Coherence in this context has a meaning somewhat different from, although related to, its usual usage in FLT. In FLT, the quasiparticles are said to propagate "coherently" in that their scattering rates are much smaller than their energies. They therefore exist as well-defined excitations. This notion of coherence is central to FLT and to its explanation of the absence of quasiparticle scattering. The absence of quasiparticle scattering is necessary to explain the successes of nearly free electron models of metals and even to explain mundane and familiar properties of metals such as their high conductivities. Coherence is thus an essential, defining property of Fermi liquids. The natural question is whether all bulk metals have coherent quasiparticles.

The existence of coherently propagating quasiparticles implies that the wave nature of the quasiparticles will be detectable in interference experiments. Exhibiting observable interference effects between different histories in this manner is the defining property of a more general notion of coherence: quantum coherence. Hereafter, we use coherence to refer to the more general notion of quantum coherence, but the reader should bear in mind the close connection between the two, especially the fact that Fermi liquid coherence implies various kinds of quantum coherence. The absence of this kind of coherence therefore implies the absence of Fermi liquid behavior. In fact, we will argue that some non-Fermi liquids can be coupled perturbatively with a single particle hopping without exhibiting coherence for the interliquid transport. In the low-energy, long-time limit, these models should have no observable interference effects between histories in which fermions move between the coupled liquids. If this occurs, then the resulting state is not a Fermi liquid.

Experimentally, we will discuss measurements of various interference effects, but on the Gedanken level, we are interested in one particular quantity that offers a natural probe of the coherence of interliquid transport. We consider two systems of interacting fermions, each initially in their ground states and identical except that the systems are at slightly different densities; that is, $\delta N = N_1 - N_2 \neq 0$, where N_1 and N_2 are the total particle numbers in each chain. At time $t = 0$, we imagine suddenly switching on the coupling between the systems in the form of a single particle hopping, $t_\perp$, and we examine the behavior of the expectation value $\langle \delta N(t) \rangle$ (*7*). It turns out that, by this measure, coupled Fermi liquids, rather remarkably, exhibit macroscopic quantum coherence (*8*), in the form of oscillations of $\langle \delta N(t) \rangle$ with vanishing damping. To better understand why undamped oscillations in $\langle \delta N(t) \rangle$ represent macroscopic quantum coherence (and why we are studying such a quantity at all), one needs to understand the close analogy between $\langle \delta N(t) \rangle$ and the quantity $\langle \sigma^z(t) \rangle$, which plays a central role in the prototypical model for the quantum-to-classical crossover, the Caldeira-Leggett model (CLM).

The CLM focuses on a two-state "macroscopic" degree of freedom described conveniently with the use of the Pauli matrices for spin one-half. This degree of freedom σ is coupled to an environment of infinitely many microscopic degrees of freedom, represented by the simplest possible realization, a bath of harmonic oscillators. The oscillators "measure" (*9*) the σ^z state of the degree of freedom, and an additional perturbation (proportional to σ^x) mixes the two σ^z states coherently. For strong enough coupling to the oscillators, the macroscopic degree of freedom should decohere, which makes superpositions of σ^z states meaningless, resulting in "classical behavior." The quantity most frequently studied is the expectation value $\langle \sigma^z(t) \rangle$, where the system has been prepared by clamping σ^z to $+1$ for all $t < 0$ and allowing the environment to adapt to this configuration. This restriction resembles an experimentally realizable situation: the controllable, macroscopic degree of freedom is held in a particular state, and the microscopic degrees of freedom, uncontrolled by the experimenter, relax to their equilibrium under these circumstances. The system is then released, and the experimenter looks for quantum interference effects in the ensuing behavior of the observable, macroscopic degree of freedom.

One can make a canonical transformation to the CLM, changing basis to the eigenstates of the joint oscillator–degree of freedom system (in the absence of the σ^x perturbation that mixes the two σ^z eigenstates). The CLM Hamiltonian then takes the form

$$H_{\mathrm{CLM}} = \frac{1}{2}\Delta(\sigma^+ e^{-i\Omega} + \sigma^- e^{i\Omega}) + \sum_i \left(\frac{1}{2} m_i \omega_i^2 x_i^2 + \frac{1}{2m_i} p_i^2\right) + \varepsilon\sigma^z \quad (1)$$

where i in the first term is -1, $\Omega = \Sigma_i (C_i / m_i \omega_i^2) p_i$, C_i is the coupling to the ith harmonic oscillator, and m_i, ω_i, x_i, and p_i are the mass, frequency, position, and momentum of the ith oscillator. For the moment, we consider only the case where the energy bias $\varepsilon = 0$. In this language, the "measurement" effects of the environment are encoded in the nondegeneracy of the states that are connected by the Δ term. In the absence of coupling to the bath (all C_i zero), the Δ term connects two degenerate σ^z states. In the presence of the bath, however, decoherence results when this degeneracy is sufficiently reduced by the operator $e^{\pm i\Omega}$, which creates and destroys oscillator bosons over a broad energy range whenever a transition between the σ^z states takes place.

In the new representation, the usual CLM preparation amounts to taking the system to be in one of the two ground states of the system in the absence of the Δ term and then suddenly switching this term on at time $t = 0$. This treatment is parallel to the preparation for coupled systems of interacting fermions; σ^z plays the role of δN, Δ the role of $t_\perp$, and the oscillator bath the role of the charge- and spin-density oscillator modes of the coupled fermions systems. In both cases, one follows the dynamics of the expectation value of a discrete, observable "macroscopic" variable (either σ^z or δN) that has been set up in a nonequilibrium state defined as a valid ground state of the problem in the absence of the perturbation (either the σ^x term or the interliquid single particle hopping term). Like the Δ term in the CLM, the action of the single particle hopping can be written, at least for Fermi and Luttinger liquids, as the product of an operator whose only action is to change δN and an exponential in the creation and annihilation operators of the charge- and spin-density oscillator modes of the interacting fermions (*10*). The resulting Hamiltonian is strikingly similar to that of the transformed CLM H_{CLM}. Although the models cannot be mapped into one another (for example, the δN variable has many values, not just two), the analogy does motivate the proposal that both coherent and incoherent dynamics can occur for δN in a manner similar to the incoherence for σ^z (*11*).

The signature of quantum coherence in the CLM is taken to be the presence of oscillations in $\langle \sigma^z(t) \rangle$, which, when present, result from interference between histories in which σ^z varies differently. We correspondingly take the presence of oscillations in $\langle \delta N(t) \rangle$ as the signature of quantum coherence for interliquid hopping.

In FLT, $\langle \delta N(t) \rangle$ exhibits oscillations with frequency $Zt_\perp$ (where Z is the overlap between the quasiparticle and free-electron wave functions at the Fermi surface) and damping that vanishes in the limit of vanishing $t_\perp$ and $\delta N(t = 0)/L$ (where L is the system size). Therefore, FLT exhibits a dramatic instance of macroscopic quantum coherence: not only are the oscillations observable, but they are essentially undamped. The source of this persistence of oscillations can be traced back to the fact that, although δN represents a macroscopic variable that might be expected to couple to a large number of uncontrolled (and potentially dephasing) microscopic degrees of freedom, FLT dictates that $\langle \delta N(t) \rangle$ is determined by the sum of many independent, coherent quasiparticle channels. Because each channel decouples from its environment in the limit of vanishing $\delta N(t = 0)/L$, coherence is unavoidable. However, for non-Fermi liquids, such as the Luttinger liquids of the coupled chains problem, there is no such special protection for the coherence of δN.

In fact, whereas FLT is analogous to the CLM with no coupling between the oscillators and σ^z, coupled Luttinger liquids are analogous to the CLM with finite coupling to an ohmic bath of oscillators (*7, 11*). As a result, the most likely behaviors for $\langle \delta N(t) \rangle$ fall into three categories: (i) for weak interactions, coherence and the characteristic oscillations, (ii) for very strong interactions, localization with $\langle \delta N(t \to \infty) \rangle \neq 0$, and (iii) for a range of intermediate interactions, incoherence, with no oscillations in $\langle \delta N(t) \rangle$ but $\langle \delta N(t \to \infty) \rangle = 0$. A remarkable result for the CLM is that the final possibility occurs over a broad range of coupling constants (*12, 13*), so that it is not the case that the oscillations simply become more heavily damped as the localization behavior is approached. Rather, the oscillation frequency actually vanishes at some intermediate coupling before the localization sets in. The physical ingredients for this behavior are present in the coupled Luttinger liquid problem (*11*), and we believe that case (iii) should occur there as well.

If the oscillation frequency of $\langle \delta N(t) \rangle$ is identically zero (over some range of couplings where $t_\perp$ is the leading instability of the uncoupled chains fixed point), then it is likely that all interchain interference effects are unobservable in the low-energy limit for this range of couplings. After all, there is nothing special about $\langle \delta N(t) \rangle$. In the CLM, it is generally believed that the lack of coherence in $\langle \sigma^z(t) \rangle$ signals a general loss of quantum coherence, and, in fact, the model was intended in part to explain the absence of interference effects for macroscopic objects. This absence is manifestly generic, rather than being limited to particular experiments. It is therefore quite likely that the disappearance of interference oscillations in $\langle \delta N(t) \rangle$ represents a generic loss of coherence, rather than one limited to the probe considered.

For two chains in the incoherent phase, we expect no pair of split Fermi surfaces, which translates, for infinitely many coupled chains, into the absence of warping of any higher dimensional Fermi surface. This lack of warping implies that this regime constitutes a new state of matter because the Fermi surface shape gives a clear zero-temperature, infinite-time distinction between the incoherent phase and a normal metal. We therefore refer to such a phase as a quantum-classical metal (QCM). We believe that this state is separated by a zero temperature quantum phase transition from a state where the interchain hopping is coherent and a three-dimensional, Fermi liquid metal occurs. There should be many differences in physical properties between these states: for example, in a QCM, the transverse electrical conductivity should lack a Drude peak, and the single particle Green's function should not exhibit a pole on the real axis that disperses with $k_\perp$ (*11*).

It is important to emphasize that our proposed phase is not one in which the electrons are confined to the chains; such a configuration would be the analog of case (ii) above for the CLM. The latter is a phase in which, in the language of the renormalization group, $t_\perp$ is an irrelevant operator. In our proposed phase, only the coherence is confined, not the electrons; diffusive interchain motion still takes place, and $t_\perp$ is a relevant operator.

Our proposal is not limited to the case of coupled Luttinger liquid chains. Although this was the best controlled case to study theoretically and the one most closely related to the CLM, there are other possibilities. In particular, a set of strongly interacting two-dimensional systems whose isolated ground states were non-Fermi liquid metals would be natural candidates for incoherent interplane hopping. It is this possibility that we believe is experimentally realized in the organic conductor $(TMTSF)_2PF_6$.

The Quantum-Classical Metallic State in $(TMTSF)_2PF_6$

The compound $(TMTSF)_2PF_6$ is a Bechgaard salt composed of linear stacks of tetramethyltetraselenafulvalene cations (Fig. 1); the stacks are arranged into planes separated by PF_6 anions, which provide overall charge neutrality and stabilize the structure. The material is highly anisotropic with a resistive anisotropy of $1{:}100{:}10^5$ at room temperature and has a single, half-filled band (*14*). It is triclinic, so that the lattice vectors are not orthogonal but, roughly speaking, the *a* axis (the most conducting direction) lies along the stacks, the *b* axis (the next most conducting direction) lies in the TMTSF planes, and the *c* axis (the least conducting direction) points out of the planes. At ambient pressure, $(TMTSF)_2PF_6$ is a spin density wave insulator, but at pressures above about 6 kbar, the ground state is superconducting. It is at such pressures and in finite magnetic fields (*15*) that we believe the incoherent interplane transport is realized. The theoretical picture (*11, 16*) is that in zero field, the interplane hopping is just barely sufficient to stabilize a three-dimensional Fermi liquid (were it not for the superconducting transition). For a magnetic field applied along the *c* axis, which minimally disrupts interplane motion, the superconductivity can be removed while retaining interplane coherence, and the behavior should be roughly like that of a Fermi liquid. For other orientations, the magnetic field interferes with the interchain coherence (*17*) by adding an effective inelasticity to the interchain hopping (*16*). This addition is analogous to the ε term in the CLM (Eq. 1); such an inelasticity reduces coherence in the CLM (*12*) and can be expected to do so here as well. In fact, data show that even moderate fields in the *b* direction completely remove interplane coherence.

The most natural experimental probes for coherence effects are low-temperature magnetotransport measurements. A number of anomalies in such measurements attracted our attention to the material. Consider the data depicted in Fig. 2 (*18*). Such an anisotropic material should have a Fermi surface consisting of a pair of well-separated sheets, and therefore, the magnetoresistance (MR) in the most conducting direction is expected in FLT to be small and to saturate quickly (*19*). Instead, the material displays an enormous, angle-dependent MR for fields rotated in the *bc* plane and current in the *a* direction. Particularly striking are the dip features that occur when the magnetic field parallels a real-space lattice direction, the so-called "magic angles" orientations. In our proposal, the dips are naturally explained as places where the magnetic field is ineffective in disrupting interplane hopping and coherence is not fully destroyed. Such a picture has a number of qualitative predictions, such as the narrowing of the *c*-direction dip roughly linearly with magnetic field, and the hierarchy of

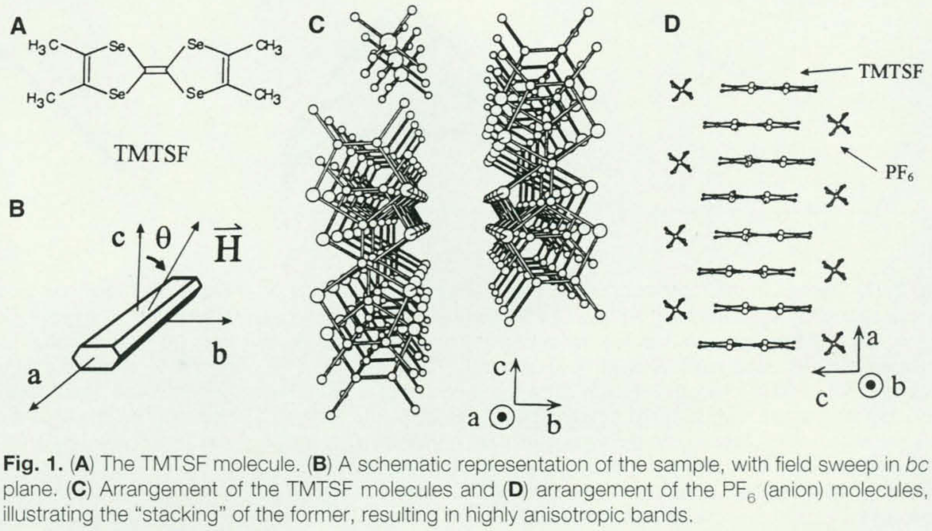

Fig. 1. (**A**) The TMTSF molecule. (**B**) A schematic representation of the sample, with field sweep in *bc* plane. (**C**) Arrangement of the TMTSF molecules and (**D**) arrangement of the PF_6 (anion) molecules, illustrating the "stacking" of the former, resulting in highly anisotropic bands.

the appearance of the dips (the c dip must appear first as a function of field strength) (11); however, it is away from the dips, where the QCM should be realized, that the strongest experimental consequences of the confinement of coherent motion to the planes should be apparent. In that state, it is impossible to observe interference effects between interplane histories, and therefore, the orbital MR contribution from the components of the magnetic field lying in the ab plane must vanish identically.

Therefore, the MR data away from the magic angles are also plotted in Fig. 2, not only versus angle (20), but also versus the component of the field perpendicular to the ab plane. Note the extent to which the data away from the magic angles collapse onto a single curve, signaling the predicted independence of the MR from in-plane field strength. The data from within the magic angle dips, where some coherence is present in our picture, are not at all independent of in-plane field strength. This deviation from scaling is a marker for the transition from incoherence to coherence; its vanishing shows the transition to the incoherent phase, expected to be truly realized only in the limit of zero temperature and a pure system, but closely approximated by the real experimental data (Fig. 2).

There is at present no other theoretical explanation for this behavior [for example, all the scenarios in (21) fail to exhibit the scaling observed in Fig. 2]. At a minimum, the data imply that away from the dips, and only away from the dips, the effect on the resistivity in this direction of interference effects due to fields in the ab plane is essentially zero. Because there might be reasons for this behavior other than the incoherence of interplane hopping, we consider further possible probes of the coherence of interplane hopping.

The most natural quantity to consider in probing c-axis coherence is the conductivity in this direction. In a simple noninteracting quasiparticle model with an isotropic scattering rate, the conductivity tensor can be calculated in the semiclassical relaxation time approximation (22)

$$\sigma_{ij} = e^2 \int \frac{d\mathbf{k}}{4\pi^3}\, \tau v_i(\mathbf{k}) \bar{v}_j(\mathbf{k}) \left(-\frac{\partial f}{\partial E}\right) \qquad (2)$$

where e is the electron charge, τ is the relaxation time, and $\bar{v}_j(\mathbf{k})$ is the weighted average of the quasiparticle velocity over a semiclassical trajectory in phase space ending at $\mathbf{k}$ (22). In the limit of large anisotropy (the smallest bandwidth negligibly small compared to the next smallest and so on), the zz component of the resistivity obtained from this approach is just $R_0[1 + (\tau l_c v_F e \mathbf{H}\cdot\hat{a} \times \hat{c})^2]$, where units are such that $\hbar c = 1$ R_0 is the zero field value, l_c is the size of a unit cell in the c direction, and v_F is the Fermi velocity along $\hat{a}$ (that is, in the most conducting direction), a well-defined quantity in the limit of large anisotropy. In low fields, we expect coherence, and our theory of the MR predicts that something like the above behavior should be observed in the c-direction resistivity, whereas in high fields, one expects the magic angle dips and the resistivity away from the dips to be independent of the in-plane field strength. In recent data (23) (Fig. 3), the crossover from the low-field, approximately FLT behavior to the magic angle behavior is striking. Again, away from the magic angle dips, the resistivity depends only on the out-of-plane component of the field, which is now the field nearly parallel to the current. At a mini-

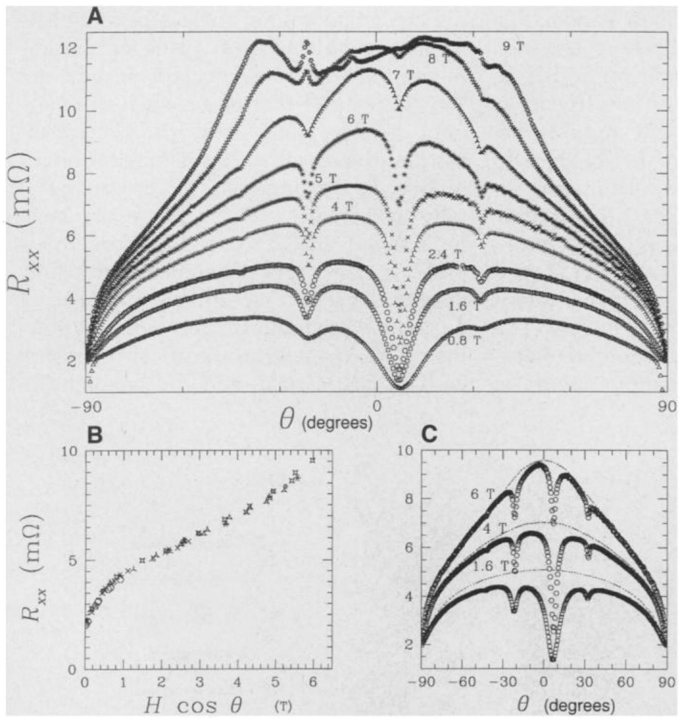

Fig. 2. (**A**) The MR in the most conducting lattice direction, a, as fields of various strengths are rotated in the plane of the other two lattice directions, that is, from $-b$ through c on to b; θ is measured from the perpendicular to the b axis (Fig. 1B). Data were taken at 10 kbar and 0.5 K. (**B**) Subset of the data from (A) for field orientations away from the magic angles. Data are replotted as resistance versus field strength out of the ab plane. Note the collapse onto a single scaling curve irrespective of field strength. (**C**) Data for field strengths of 1.6, 4, and 6 T together with the expected MR from the scaling curve of (B). Deviations from scaling occur only within the vicinity of the magic angles that decrease rapidly with field.

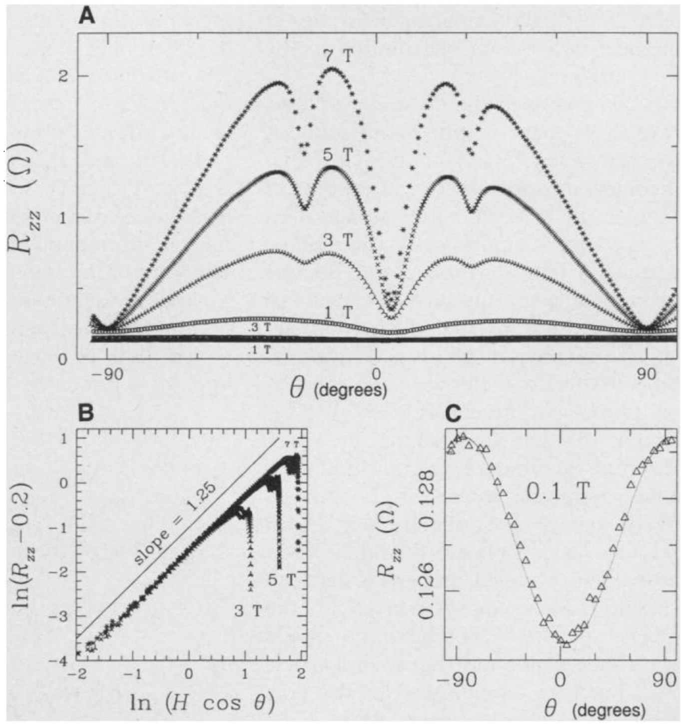

Fig. 3. (**A**) The MR perpendicular to the ab plane as fields of various strengths are rotated as in Fig. 1. Data were taken at 10 kbar and 1.3 K (26). (**B**) Data from (A) for 3, 5, and 7 T plotted as the natural logarithm of deviation from a reference value versus the natural logarithm of magnetic field strength perpendicular to the ab plane. Away from the magic angles, the data exhibit dependence on only one component of the magnetic field. The dependence is well described by $\Delta R \propto (H \cos \theta)^p$, where $p \sim 1.25$. This nontrivial power law is not compatible with existing FLT descriptions of MR. (**C**) Weak-field (0.1 T) MR for bc rotation. The dotted line is a fit to the data of the form $R_0 + \alpha|H \times \hat{c}|^2$, the semiclassical prediction for the limit of extreme anisotropy.

mum, the data require that away from the dips, and only away from the dips, the sum of the interference effects attributable to fields in the *ab* plane on the resistivity out of the *ab* plane is nearly zero. Again, there is no theoretical proposal other than the incoherence of the interplane hopping that can account for the observed effect, especially given the qualitative agreement between the low-field data and the expected behavior of a relatively clean Fermi liquid.

Danner and Chaikin (*5*) have examined another interference effect (*24*) that is sensitive to the coherence of *c*-axis transport and the existence of a three-dimensional Fermi surface in $(TMTSF)_2PF_6$. The results are again naturally explained by the presence of coherence in fields whose projection along certain real space lattice vectors is sufficiently small, but the total absence of interplane coherence for other fields.

It is, of course, impossible to demonstrate the total absence of coherence experimentally: one can only show that various measured quantities are consistent with the absence of interplane coherence, that is, they exhibit no signs of interference effects between histories involving interplane motion. It is remarkable, however, that three different measures of interplane coherence show the complete absence of such interference effects. The MR out of the *ab* plane is particularly striking, because it most directly probes the *c*-axis charge transport and, although its low-field behavior fits well into the expectations of FLT, its behavior in high fields is totally anomalous from the FLT point of view.

If one accepts that there is no coherent interplane motion, there is no other explanation besides our proposal of relevant interplane hopping that has been driven incoherent by in-plane interaction effects. If the hopping were irrelevant—that is, if the effective low-energy theory describing the system had no out-of-plane hopping—the unimportance of magnetic fields in the *ab* plane would follow naturally. What would not make sense, however, would be the strong angular dependence of the *c*-direction MR (particularly the dips) or even the mere existence of a substantial low-temperature *c*-axis conductivity. Further, if the low-energy theory had no *c*-axis hopping of charge carriers, then the resistivity would have to diverge as $T \rightarrow 0$. The experimental behavior in the incoherent phase (Fig. 4) yields no evidence for such behavior down to below 1 K.

Moreover, measurements at 50 mK (Fig. 4) show no signs of a diverging resistivity. Therefore, an explanation of the unimportance of magnetic fields in the *ab* plane based on the absence of *c*-axis hopping is untenable. If the hopping exists and is incoherent, it could be a property of a low-energy fixed point or it could be the result of inelastic scattering or disorder. However, because the essential features of the MR remain down to at least 50 mK (Fig. 4), an explanation based on inelastic effects again appears untenable. For many reasons, inelastic disorder scattering is also unable to account for the loss of coherence (*25*). For example, (i) the magic angles and scaling are only observed in the highest quality crystals; (ii) the low-field data are consistent with a FLT type behavior with a scattering rate of about 1 K, a rate too low to credibly explain the incoherence of *c*-axis transport; and (iii) the magic angle dips in themselves demonstrate a strong dependence of the in-plane transport on the coherence of the *c*-axis motion, a dependence that is impossible if disorder dominates *c*-axis motion.

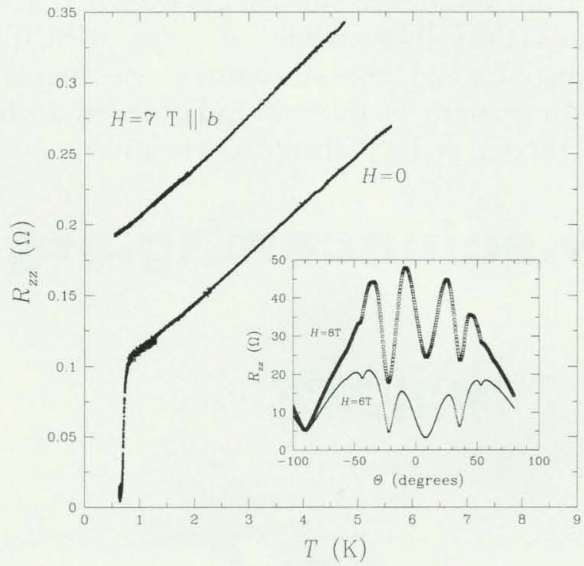

Fig. 4. Temperature dependence of the resistivity perpendicular to the *ab* plane both in the absence of a magnetic field and in a field of 7 T along *b* (the latter condition is in the incoherent phase). Data taken at an applied pressure of 10 kbar. (**Inset**) MR perpendicular to the *ab* plane for *bc* field rotations as in Fig. 2, but at 8.2 kbar and 50 mK (*27*).

Conclusion

The experimental situation is therefore remarkably compelling. All indications are that in the limit of a pure system and zero temperature, $(TMTSF)_2PF_6$ exhibits a phase in which an applied magnetic field (*15*) not only destroys superconductivity but drives the system to a state of matter characterized by finite, relevant interplane electron hopping but the complete absence of observable interference effects between histories involving interplane motion. More succinctly, at zero temperature, there is a nonvanishing interplanar electron conductivity, but it is completely incoherent, in contrast to the coherent in-plane transport. The state is, therefore, a new phase of electronic matter, a non-Fermi liquid metal characterized by "quantum" in-plane and "classical" interplane transport—in short, a quantum-classical metal.

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25. Disorder in the hopping between planes is an even less tenable proposal than in-plane disorder; as long as there is a spatially uniform component to $t_{\perp}(x)$, its Fourier transform has a delta function at zero momentum transfer with finite weight. This finite amplitude for momentum-conserving interplane motion of a single electron is generically far more relevant than any momentum-nonconserving hopping process that might be invoked to account for the loss of interplane coherence and should always dominate over such processes in the low-energy limit.
26. The experiments were performed in a split-bore superconducting magnet using a ^{3}He refrigerator capable of cooling to 0.5 K. The small dimensions of the pressure cell (1.39 cm by 0.76 cm) [see G. M. Danner and P. M. Chaikin, *Rev. Sci. Instrum.* **66**, 3951 (1995)] allowed us to rotate the sample in two mutually perpendicular planes by 4π steradians and thus align the crystal axes with respect to the magnetic field to better than 0.05°. The sample typically had a needle shape with size 1 mm by 0.15 mm by 0.1 mm. Resistance measurements were made with a standard four-probe low-frequency ac technique. The magnitude and anisotropy of the resistance, MR, and critical field were consistent with previous measurements and indicate the high quality and extremely long mean free paths that characterize this class of materials.
27. The low-temperature ($T \sim 50$ mK) data represented in the inset of Fig. 4 were obtained at the National High Magnetic Field Laboratory in Tallahassee, FL, using a dilution refrigerator. The experimental details are as in (*26*).
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