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AUTHOR(S): C. J. Pethick* and David Pines**
University of Illinois at Urbana-Champaign, Urbana, IL 61801

*Also at Nordita, Blegdamsvej 17, DK-2100 Copenhagen Ø, Denmark

**Also at Center for Materials Science, Los Alamos National Laboratory, Los Alamos, NM 87545

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Los Alamos Los Alamos National Laboratory
Los Alamos, New Mexico 87545

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HEAVY ELECTRON SUPERCONDUCTIVITY: FROM 1K TO 90K TO ?

C. J. Pethick* and David Pines**

Department of Physics, University of Illinois at Urbana-Champaign
1110 West Green Street
Urbana, IL 61801, U.S.A.

INTRODUCTION

Heavy electron systems are intermetallic compounds containing elements with unfilled f-electron shells, such as U or Ce, which at room temperature and above behave like a weakly interacting collection of f-electron moments and conduction electrons with ordinary masses, while at low temperatures the conduction electron specific heat becomes typically some hundred times larger than that found in most metals.[1] These highly correlated low temperature states display remarkable behavior whether the system remains normal down to the lowest temperature measured, becomes antiferromagnetic, or becomes superconducting. While in ordinary metallic superconductors a dilute concentration of magnetic impurities destroys superconductivity, in heavy electron systems superconductivity and antiferromagnetism can coexist; a transition to either ordered state may be followed by a second transition to a phase containing both states. Thus in both UPt_3 and URu_2Si_2 one finds on lowering the temperature that an antiferromagnetic transition is followed by a transition to the superconducting state, while in $U_{0.97}Th_{0.03}Be_{13}$ the order of the transitions is reversed.

In this talk we shall review the experimental results and physical arguments which led us to conclude that in heavy electron systems the physical mechanism responsible for superconductivity is an attractive interaction between the heavy electrons which results from the virtual exchange of antiferromagnetic f-electron moment fluctuations.[2] In these systems, then, the superconductivity is of purely electronic origin; the phonon-induced interaction between electrons which leads to superconductivity in ordinary metals plays little or no role.

From the perspective of scientists searching for high temperature superconducting materials heavy electron systems thus provide both good news and bad news. The good news is a purely electronic mechanism for

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*Also at Nordita, Blegdamsvej 17, DK-2100 Copenhagen Ø, Denmark

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superconductivity has been discovered; hence the "phonon-barrier," the existence of a maximum superconducting transition temperature of ~30K for metals in which electron-phonon interactions are responsible for superconductivity,[3] is broken; the bad news is that although the superconductivity is of non-phononic origin, the superconducting transition typically is found only at $T \lesssim 1K$.

Does this always have to be the case? In the next part of this talk we consider the possibility that the high temperature superconductors are indeed part of the heavy electron family, albeit a collateral branch in which magnetic excitations are present but the carrier density is so low that the screening of magnetic interactions by itinerant electrons or holes is negligible. Under these circumstances the coupling of charge carriers to spin fluctuations could easily give rise to substantial carrier effective masses of the size ($\sim 10\ m_e$) recently measured,[4] while the characteristic temperature for the spin-fluctuation induced superconductivity would be far closer to 100K than 1K.

One consequence of a large effective carrier mass is a small coherence length, $\xi_o \sim 8A$; this in turn makes possible intrinsic pinning of the magnetic vortices in the 90K superconductors by the barium or rare earth atoms which lie outside the copper-oxide planes, a possibility we consider briefly in the latter part of our presentation.

PHYSICAL PICTURE

At high temperatures heavy electron systems behave like a collection of weakly interacting f-electron moments and conduction electrons, while at very low temperatures, so far as thermal and transport processes are concerned, they behave like a system of strongly interacting itinerant electrons which scatter against impurities, against low frequency f-electron spin fluctuations, and against one another.

A physical picture of the transition between these two regimes is that as the temperature is lowered the local moments and conduction electrons become more and more strongly coupled. The magnetic behaviour is quenched while the effective mass of the itinerant electrons becomes substantially enhanced. As a consequence of this interaction, the f-electrons are no longer confined to the magnetic sites, but can hop into the conduction band, as in the Anderson model. The itinerant heavy electron states at low temperatures are therefore superpositions of localized f-electrons and conduction electrons. Their quite strong interaction reflects not so much their direct Coulomb interaction, as it does an interaction induced by their coupling to spin fluctuations on the magnetic sites, and it provides a natural explanation for the large finite temperature corrections to the low-temperature form of the specific heat, the strong temperature dependence of the electrical resistivity and other transport coefficients, and the appearance of superconductivity.

In the very low temperature limit the thermal and transport properties of heavy fermion systems in the normal state should be those expected for heavy electron Fermi liquids. However, in most cases experiments have not yet been carried out in the Landau limit, that is at temperatures sufficiently low that one can neglect, in first approximation, the frequency dependence of the quasiparticle energies and quasiparticle scattering amplitudes associated with the coupling of the conduction electrons to the localized f-electrons. If we define θ_{coh} as the temperature below which resistivities fall off sharply with decreasing temperature, then it is only at temperatures $T \ll \theta_{coh}$, that one expects to observe the Landau temperature dependence of the electrical resistivity, ρ, the thermal resistivity, W, and the ultrasonic attenuation coefficient α, in which the finite temperature corrections to the low temperature limiting behaviour are proportional to T^2. Such Landau limiting behaviour is observed for UPt_3 at temperatures below ~1.5K, while UBe_{13} at zero pressure becomes a superconductor well before it reaches a temperature at which Landau theory would apply.[5]

The strong coupling between the f-electrons and the conduction electrons which is responsible for the heavy itinerant quasiparticles gives rise to a compensating electron cloud which alters the magnetic response of the local moments. If magnetization were a conserved quantity then local moments and their corresponding electron clouds would not contribute to the long wavelength magnetic susceptibility $\chi(T)$ at low temperatures; that quantity would be entirely determined by the heavy electron quasiparticle contribution χ_{qp}. Because magnetization is not conserved there can be a significant non-quasiparticle contribution χ_{loc}, to $\chi(T)$, which arises from the polarization of the local moments and their compensating clouds, that is from virtual excitations at finite frequencies.

To see how this comes about, consider the exact expression for the magnetic susceptibility at zero temperature,

$$\chi^{M}(\underset{\sim}{q},\omega) = \sum_{n} \frac{|(M^{+}_{\underset{\sim}{q}})_{no}|^2 \, 2\omega_{no}}{\omega^2_{no} - (\omega+i\eta)^2} \tag{1}$$

where 0 denotes the ground state, n an excited state, ω_{no} the excitation energy of the state n with respect to the ground state, and $M_{\underset{\sim}{q}}$ is the magnetic moment operator. At long wavelengths the excited states may be divided into two classes: i) States obtained by destroying a quasiparticle below the Fermi surface and creating a quasiparticle just above the Fermi surface in the same band. These quasiparticle-quasihole pair states have an energy of order $v_F q$, where v_F is the Fermi velocity. ii) All other states, such as one containing a quasihole in one band and a quasiparticle in another band (an interband transition), a state containing two or more quasiparticle-quasihole pairs, or, in the case of Kondo and similar systems, a state obtained by polarizing a localized spin and its compensating electron cloud. χ^M therefore takes the form

$$\chi^{M} = \chi^{M}_{Landau} + \chi^{M}_{loc}, \tag{2}$$

where the first term comes from the single pair states (i), and the second from states (ii). By performing neutron scattering experiments at small $\underset{\sim}{q}$, one can in principle distinguish between these two contributions, since the frequencies associated with the Landau contribution all vanish for small $\underset{\sim}{q}$.

If magnetization is conserved, χ_{loc} vanishes at long wavelengths. This may be seen from the fact that the total magnetization commutes with the Hamiltonian,

$$\left[\underset{\sim}{H}, M_{q=o}\right]_{no} = \omega_{no}\left[M_{q=o}\right]_{no} = 0, \tag{3}$$

and therefore, assuming $\omega_{no}(M_{\underset{\sim}{q}})_{no}$ for $\underset{\sim}{q} \to 0$ tends to its value at $\underset{\sim}{q} = 0$, it is easy to see that χ^{M}_{loc} must vanish. This shows that for such systems the Landau contribution to the susceptibility at long wavelengths must be the total susceptibility. This is true for liquid ^{3}He: in this case the magnetization is proportional to the spin, which is conserved if one neglects the nuclear dipole-dipole interaction. Conservation of magnetization also leads to the conclusion that the magnetic moment associated with a quasiparticle is equal to the bare moment.

In heavy fermion systems, magnetization is not conserved, due to the existence of both spin and orbital contributions to it, and to spin-orbit coupling. Consequently χ^{M}_{loc} is finite in the limit $\underset{\sim}{q} \to 0$. The absence of magnetization conservation also means that the effective magnetic moment of a quasiparticle is not related in a simple way to the bare moments of either an f-electron or a conduction electron.

Inelastic neutron scattering experiments give no evidence for a component of Im χ^{M}, the magnetic structure factor, whose frequency tends to zero as $\underset{\sim}{q} \to 0$. This region is difficult to investigate directly, but the fact that the contribution to χ^{M} from the frequencies which are accessible experimentally can account for all of the measured long-wavelength susceptibility to within experimental accuracy suggests that χ^{M}_{Landau} cannot contribute more than 10-20% of the total.[6]

A PHENOMENOLOGICAL DESCRIPTION OF HEAVY FERMION BEHAVIOUR

Recently it has been shown that neutron scattering results for UPt_3, $CeCu_3$, and U_2Zn_{17} may be fit by a model for the spin-spin correlation function in which fluctuations of the magnetic moment at the f-atom site are coupled to those at other sites by an effective exchange interaction.[7] If one assumes that all the magnetic moment is associated with electrons in f-orbitals, this leads to an expression for the wavenumber- and frequency-dependent spin-spin correlation function of the form

$$\chi(\underset{\sim}{q},\omega) = \frac{\chi_{\mu}(\omega,T)}{1-J(\underset{\sim}{q},\omega,T)\,\chi_{\mu}(\omega,T)} \tag{4}$$

where $\chi_{\mu}(\omega,T)$ describes the correlations of the spin at a single f-site, including the effects of interaction with the compensating electron cloud, and $J(\underset{\sim}{q},\omega,T)$ is an effective exchange interaction which describes the coupling between spins at different sites. (For simplicity we restrict ourselves to the case where the sites of all magnetic ions are equivalent.)

In the fits to the data, J was taken to be a temperature-dependent nearest neighbour interaction, and χ_μ was taken to be of the form

$$\chi_\mu = \frac{\chi_o \Gamma}{\Gamma - i\omega} \quad (5)$$

which is known to give a good description of the properties of a single Kondo impurity. Here χ_o is the susceptibility of a single ion and Γ is a measure of the typical excitation energies for a single ion and its screening cloud, energies which lie between 50K and 250K for the systems thus far studied; for an isolated impurity, Γ would be of order the Kondo temperature T_K.

We have proposed that an expression of the form (4) provides a useful starting point for the examination of all aspects of heavy fermion behavior.[2] From a microscopic point of view the induced spin-spin interaction is given by an expression of the form

$$J(\underset{\sim}{q},\omega,T) = -\sum_{\underset{\sim}{K}_n} |V_{\underset{\sim}{q}+\underset{\sim}{K}_n}|^2 \chi_c(\underset{\sim}{q}+\underset{\sim}{K}_n,\omega,T) \quad (6)$$

where V describes the coupling of a conduction electron-hole pair to the local spin fluctuations described by $\chi_\mu(\omega,T)$, K_n is a reciprocal lattice vector, and χ_c is the conduction electron-hole spin-spin response function. As a result of the coupling, J, the characteristic energies which enter into the low frequency limit of χ, (Eq. 4), become wavevector- and temperature-dependent, being given by

$$\theta_{loc}(\underset{\sim}{q},T) \cong \Gamma[1-J(\underset{\sim}{q},0,T)\ \chi_\mu(0,T)] \quad (7)$$

We argued that the presence of a second energy scale, lower than T_K, is a characteristic feature of all heavy fermion systems, and may be a necessary condition for observing heavy fermion behavior. Put another way, if $J\chi_\mu \ll 1$, one is likely in a weak coupling limit, and no heavy fermion behavior results. One the other hand if, as in U_2Zn_{17}, for some wavevector $\underset{\sim}{q}$, and temperature T, $J\chi_\mu = 1$, then an antiferromagnetic phase transition occurs. (Indeed, Broholm et al.[7] have shown that this transition is driven by a temperature dependent coupling, J, which below 18K increases with decreasing temperature until it drives the antiferromagnetic transition at 9.7K.) Normal and superconducting heavy fermion compounds would seem to lie in the strong coupling regime, $J\chi_\mu \sim 1$.

The coupling between the heavy quasiparticle pairs and the local spin fluctuations gives rise to an induced wavevector-, frequency- and temperature-dependent, heavy electron interaction,

$$U_{ind}(\underset{\sim}{q},\omega,T) = -V_{eff}^2\ \chi(\underset{\sim}{q},\omega,T) = \frac{V_{eff}^2\ \chi_\mu(\omega,T)}{1-J(\underset{\sim}{q},\omega,T)\ \chi_\mu(\omega,T)} \quad (8)$$

The matrix element, V_{eff}, includes vertex corrections to the electron-local momentum dependence of $U(\mathbf{q},\omega,T)$ would arise from that of $J(\mathbf{q},\omega,T)$. For frequencies low compared to the characteristic frequencies which enter into χ, that interaction will be attractive between like spins and repulsive between unlike spins; to the extent that χ exhibits antiferromagnetic correlations (and neutron scattering experiments suggest that this might quite generally be the case), $U_{int}(q,\omega)$ will behave in similar fashion. This induced interaction is the physical origin of both the $T^3 \ln T$ corrections to the specific heat (where these are observed) and of the pairing instability which gives rise to superconductivity. The proposed approach is quite reminiscent of the electron-phonon interaction problem, with the local moment spin fluctuation frequency-dependent susceptibility playing the role of a phonon propagator. However, there is no reason to expect that a Migdal theorem exists for the heavy electron local moment fluctuation interaction. Indeed, in the present theory there is a considerable amount of feed-back, and possible non-linear behavior, in that, for example, J depends on χ_c which in turn depends on J through electron-local moment fluctuation coupling.

Transport coefficients depend on the behavior of J at large wavevectors, since scattering phenomena are dominated by the coupling of heavy electrons to large wavevector moment fluctuations. A test of this hypothesis, and of the overall model, is obtained by examining the changes in the resistivity as a function of pressure and magnetic field, in an approach which attributes such changes either to changes in θ_{loc} which proceed à la Kondo, $[\theta^2_{loc}(H) = \theta^2_{loc} + \mu^2_{loc} H^2]$, and/or to changes in J. Batlogg[9] has recently found that such scaling arguments work quite well for the resistivity and magnetization of UBe_{13} in quite large magnetic fields.

Finally, we argued that the mass enhancement of heavy fermions arises from their coupling to local moment fluctuations. As was the case for transport phenomena, the calculated state density will depend primarily on the coupling of the conduction electrons to the large wavevector local moment fluctuations.

Recently Norman[9] has carried out a microscopic model calculation for UPt_3 which serves as a test of our proposed phenomenological approach. He has assumed that quasiparticles on a Fermi surface which is consistent with the de-Haas van Alphen results of Taillefer et al.[10] are coupled to moment fluctuations whose spectra are those measured by Aeppli et al.[7]; he finds both a mass enhancement and a superconducting transition temperature which are in good qualitative agreement with experiment.

Our physical picture and phenomenological description may be rich enough to make possible an understanding of the extraordinary, and diverse, sensitivity of various heavy fermion physical phenomena to pressure and to the presence of impurities. For example, impurities can alter χ by changing either χ_μ and/or J. Either, or both, of these quantities may in turn be quite sensitive to changes in density; moreover, the introduction of impurities can give rise to local changes in density. As a result, a natural explanation may emerge for the fact that in heavy fermion systems the thermal expansion, most often negative, is some four orders of magnitude larger than that of an ordinary metal; the observed values of magnetostriction exceed those of transition metals by two or more orders of magnitude, and, finally, the introduction of impurities can bring about changes in the resistivity which can be two orders of magnitude greater than the value obtained from an estimate based on using for the quasiparticle-impurity scattering amplitude the unitarity limit in a single partial wave.

HEAVY FERMION SUPERCONDUCTIVITY

A fundamental question concerning superconductivity in heavy electron systems is whether it is the heavy electrons that become superconducting. Clear evidence for the pairing of the heavy electrons is provided by measurements which show that the jumps in the specific heat at the transition temperature, T_c, to the superconducting phase, are comparable to the specific heat in the normal phase.

A second fundamental question is whether the superconducting energy gap has nodes on the Fermi surface, and, if so, what their character is. Experimentally, no equilibrium or transport properties in the heavy fermion superconductors exhibit the exponential behavior expected for states with a non-zero energy gap everywhere on the Fermi surface; rather both specific heat and transport measurements display the power-law behavior characteristic of states with gaps which vanish at points or along lines on the Fermi surface. Specific heat measurements at low temperature, which reflect the density of quasiparticle states at energies of order k_BT, give direct evidence about the nodes of the gap. At low temperature, the only quasiparticles excited will be those in the vicinity of nodes of the gap. These states possess an energy less than k_BT and lie within an angle $\sim T/\Delta$ of a node, where Δ is the maximum value of the energy gap on the Fermi surface. A simple geometric argument shows that the density of quasiparticles varies as T^2 for nodes at points and as T for nodes on lines, and the corresponding variation of the specific heat is as T^3 and T^2, respectively. In this way the experimental measurement of a T^2-dependence of the specific heat for UPt_3 shows that the energy gap vanishes on a line or lines, while the T^3 dependence found in UBe_{13} is indicative of a gap which vanishes at points. Thus heavy-fermion systems possess at least two superconducting states. Since UBe_{13} possesses cubic symmetry, while UPt_3 is hexagonal, it is possible that crystal structure plays a role in determining the nature of the superconducting state. Evidence that suggests the possible existence of two superconducting states in a single system is provided by specific heat and critical field experiments on $U_{1-x}Th_xBe_{13}$, where x, the concentration of Th impurities, lies between 2 and 4 percent.

A third question of interest is where the nodes lie on the Fermi surface. Information about this is contained in measurements of transport coefficients such as accustic attenuation. In UPt_3 the attenuation, α, of transverse ultrasound propagating in the basal plane, measured by Shivaram et al.,[11] shows a different temperature dependence according to whether the sound wave is polarized in the basal plane ($\alpha\propto T$) or perpendicular to it ($\alpha\propto T^2$). These results suggest that quasiparticles move more freely in the basal plane than perpendicular to it, which would be consistent with a quasiparticle gap having nodes on lines on the Fermi surface perpendicular to the hexagonal axis. Further evidence for this behavior of the gap is provided by the recent tunneling measurements of Batlogg et al.,[12] which give no evidence for a gap when quasiparticles are injected across crystal faces with normals perpendicular to the hexagonal axis, but show a distinct gap when quasiparticles are injected across faces with moments parallel to the hexagonal axis.

A considerable amount of effort has gone into trying to understand transport in the superconducting states. Under circumstances in which scattering by impurities is the dominant process, as is the case in UPt_3 at temperatures of the order of T_c and lower, the temperature dependence of the transport coefficients seems to disagree with calculations for any anisotropic superfluid state if the scattering is treated in the Born approximation. In this approximation the lowest order s wave scattering by a single impurity is considered; the calculated mean free paths increase with decreasing temperature, and one finds results for the thermal conductivity, κ, and acoustic attenuation, α, which are much larger than those observed experimentally. We have shown[13] that if one takes into account the multiple scattering of quasiparticles by impurities, and if one is near the unitarity limit characterized by a phase shift, $\delta \sim \pi/2$, the mean free path for electron-impurity scattering shows remarkably little dependence on temperature, so that both α and κ/T fall off with decreasing temperature, in agreement with experiment. The transport data for UPt_3, including the anisotropies observed by Shivaram et al.[11] in the attenuation of transverse sound, can be accounted for qualitatively if, as noted above, one has a polar state in which the superconducting gap has nodes on lines on the Fermi surface which are parallel to the c axis of the crystal, and the mean free path is independent of temperature.[14] In our calculations, we did not take pair-breaking into account. Pair-breaking effects are important only at energies close to the gap energy, Δ, and at low energies, $E \sim \hbar/\tau_N$, where τ_N is the lifetime for impurity scattering in the normal state; these have been included in the work of Schmitt-Rink et al.,[14] Hirschfeld et al.,[15] and Scharnberg et al.[16] who find in numerical calculations that with $\hbar/(\tau_N \Delta) \sim 10^{-2}$, pair-breaking effects are important for polar states only at temperatures below $\sim(T_c/10)$, in agreement with the above estimate.

Quite generally features around the nodes are smeared out by impurity scattering. Evidence for this physical affect on the density of states in the superconducting state of UBe_{13} has been found by Ott et al.[17] in experimental measurements of the specific heat at low temperatures ($T \gtrsim 50mK$); the experimental results are in excellent agreement with theoretical calculations of the state density which assume an axial state, in which the energy gap has point nodes, and electron impurity scattering which is near the unitarity limit.

There can be little doubt that the superconducting states observed in the heavy electron systems are unconventional, when compared to typical metallic superconductors. While there is as yet no theoretical proof or direct experimental demonstration that electron-phonon interactions are essentially irrelevant to heavy fermion superconductivity, in view of the persuasive physical arguments that the physical origin of the large masses is the coupling of conduction electrons to the local moment fluctuations, and that the virtual exchange of such spin fluctuations gives rise to an attractive interaction between heavy electron quasiparticles, and the model calculation of Norman,[9] it would seem overwhelmingly likely that it is the electron-local moment fluctuation coupling which is responsible for heavy electron superconductivity. Whether the resulting pairing state is "p-like" or "d-like" depends on the details of the wavevector dependence of the effective attractive interaction, and present evidence clearly favors the latter possibility.

ARE THE HIGH TEMPERATURE SUPERCONDUCTORS A BRANCH OF THE HEAVY ELECTRON FAMILY?

Since heavy electron systems provide us with a new mechanism for superconductivity and new pairing states in metals, it is natural to inquire whether the physical origin of superconductivity in the ceramic oxides is similar; thus does superconductivity in these systems arise from an attractive interaction between electrons or holes induced by their coupling to spin fluctuations excitations, and are these superconductors another branch of the heavy electron family? Models in which spin fluctuations induce superconductivity have been proposed by a number of authors,[18] while the detection of antiferromagnetic ordering in La_2CuO_{4-y} for non-zero values of y,[19] provides support for this hypothesis. The very recent measurement[4] of itinerant carrier masses of the order of ten times the bare electron mass in $YBa_2Cu_3O_7$ makes the family resemblance to heavy electron systems still more striking. Quite generally in the high T_c superconductors, the copper-oxide planes represent a promising source of spin fluctuations, while if the itinerant carriers belong to a distinct, but nearby (in energy) band, the basic physics would be remarkably similar, with a spin-fluctuation induced interaction being responsible for both the heavy carrier mass and superconductivity; it is also possible that the carriers, which are hole-like for $YBa_2Cu_3O_7$,[4] and spin fluctuations are excitations which belong to the same band.

The reason, then, that one achieves high superconducting temperatures in the ceramic oxides is that the itinerant carrier density is so small that the induced spin-spin interaction, J, of Eqtn. (6), plays almost no role; rather the magnetic behavior is associated with exchange interactions between spins on nearest neighbor and next nearest neighbor sites, and their scale is set by the Neel temperature for the copper-oxide layers; since this temperature can be of the order of room temperature, superconductivity at high temperatures can easily result. Put another way, what spoils the chances for high T_c in the heavy electron systems is that the conduction electrons are sufficiently dense to screen the f-electron moments, so that the scale over which the heavy electron interaction can be attractive is one or two orders of magnitude smaller than the Curie-Weiss temperature; that same phenomenon is responsible for the fact that the Neel temperatures in heavy electron systems are $\lesssim$ 20K, rather than being comparable to or greater than room temperature. Since spin-orbit coupling effects are small in the ceramic oxides, the itinerant carrier long wavelength magnetic susceptibility will be of the Pauli-Landau type, and will contain no local moment contributions.

It is worth remarking that the measured low itinerant carrier densities in the ceramic oxides ($n \sim 3 \times 10^{21}$ cm^{-3}) correspond to values of $r_s \sim 8$, where r_s is the interelectron spacing divided by the Bohr radius. As r_s increases beyond its values at ordinary metallic densities ($r_s \lesssim 3$) the Fermi liquid parameter F_0^a, associated with direct Coulomb correlations, becomes increasingly negative[20]; thus at $r_s \sim 8$, direct Coulomb correlations favor both the inferred substantial exchange enhancement of the Pauli magnetic susceptibility,[21] and p state or d-state superconductivity, depending on the wavevector dependence of that static susceptibility.[2,22]

It is possible therefore that direct Coulomb correlations act to enhance the spin-fluctuation-induced interactions associated with the copper-oxide layers, and so further increase the superconducting transition temperatures.

In common with the plasmon-exchange and exciton-exchange mechanisms, the spin-fluctuation mechanism provides a natural explanation for the observed absence of an isotope effect in $YBa_2Cu_3O_7$ and $EuBa_2Cu_3O_7$.(23) Unlike the above mechanisms, it also provides a natural explanation for the extreme sensitivity of the copper-oxide superconductors to substitutions for the copper ions. Such substitutions, it may be argued, can change dramatically the nature of the spin fluctuation excitation in the copper-oxide layers, and easily destroy superconductivity, while it is difficult to see why these substitutions would affect the exchange of virtual plasmons or excitons between carriers, and hence affect any superconductivity arising from that exchange. Spin-fluctuation exchange mechanisms will tend to give rise to energy gaps with nodes, and at present there is no direct evidence for such gap behavior.(24) It is possible that it is there, but masked by the influence of anisotropy and/or scattering; if so, transport and specific heat experiments on single crystals may be expected to probe gap structure, if any exists. Finally, we note that for spin-fluctuation mechanisms, the maximum temperature for a superconducting transition will be of order the Neel temperature, T_N, as has been noted by deGennes;(18) materials with an underlying large Neel temperature, or large energy spin fluctuations, and low carrier concentrations would appear to be promising candidates for superconductors with transition temperatures well in excess of 90K; hence the question-mark in our title.

Much theoretical and experimental work will be required to test the spin fluctuation mechanism for ceramic oxide superconductivity. To cite but two examples, inelastic neutron scattering experiments on single crystals will test whether the superconducting materials possess spin fluctuation excitations of the desired character, while the two-dimensional character of the layers, and anisotropic effects more generally, may well play a special role. It took some three years of intensive experimental and theoretical investigations for the heavy electron community to arrive at a consensus on the physical picture we have set forth in this article; it would not be surprising if a comparable period of time might be required to arrive at a comparable consensus on the new high T_c materials.

INTRINSIC FLUX PINNING IN $YBa_2Cu_3O_7$

The very short coherence length, $\xi_0 \lesssim 8Å$ which Bedell et al.(25) infer from their self-consistent analysis of experiments on both the normal and superconducting properties of $YBa_2Cu_3O_7$, opens up the interesting possibility of intrinsic flux pinning in this material, i.e. the pinning of magnetic vortices to atoms in the unit cell, rather than the extrinsic pinning to crystalline imperfections usually found in Type II superconductors. Thus it appears energetically favorable for magnetic vortices to pass through the Y and Ba atoms in the unit cell, and hence to be pinned to these atoms. The situation resembles that found in the "other" high temperature superconductors - neutron stars - in which the pinning of vortices in the rotating neutron superfluid to crustal nuclei has been shown to explain glitches and post-glitch behavior in pulsars.(26) Whether flux pinning in the terrestrial high temperature superconductors will correspond to the weak pinning or super-weak pinning situations encountered in neutron stars remains to be determined. It would seem, however, that intrinsic flux pinning both provides a natural explanation for the pinning effects observed by Harshman et al.[24] and makes possible very substantial critical currents in directions parallel to the copper-oxide planes in single crystal defect-free $YBa_2Cu_3O_7$. One would expect that intrinsic pinning phenomena will be highly anisotropic, and there is the further intriguing possibility that although the superconductivity of the 90K superconductors is not affected by the substitution of various rare earth impurities for Y, the resulting pinning phenomena and critical currents might be substantially influenced.

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