1 Introduction

The C$_{60}$ molecule is the basis for a host of interesting electronic materials. Alkali-intercalated C$_{60}$ compounds of the form A$_3$C$_{60}$ display superconductivity with transition temperatures T$_c$ exceeding 30K. The surprisingly high T$_c$ suggests either a fortunate combination of modest density of states, high-frequency phonons, and good electron coupling, or a novel superconductivity mechanism. Normal state transport measurements are a reliable method for determining relevant model parameters. Polymerized A$_1$C$_{60}$ materials have a linear-chain-like structure suggestive of anisotropic conduction. Because of the lower dimensionality, these materials may also be susceptible to density wave or other instabilities at lower temperature. Transport measurements, in conjunction with other measurements such as structural, magnetic, photoemission, and spin resonance, are again essential tools in characterizing the electronic states of these unusual polymers.

Although a wide variety of alkali-fullerene has been synthesized, few materials have been extensively characterized by transport measurements. In the case of three-dimensional alkali-fullerides, K$_3$C$_{60}$ and Rb$_3$C$_{60}$ have been the most carefully studied. In this report we shall focus on dc transport measurements on single crystal K$_3$C$_{60}$ and Rb$_3$C$_{60}$. For the polymerized materials, we review dc transport in K$_1$C$_{60}$, Rb$_1$C$_{60}$, and Cs$_1$C$_{60}$.

2 Transport in K$_3$C$_{60}$ and Rb$_3$C$_{60}$ Crystals

2.1 Dc Resistivity

Fig. 1 shows the normalized dc electrical resistivity of single crystal specimens of K$_3$C$_{60}$ and Rb$_3$C$_{60}$ below 250K. Although the functional
form of \( \rho(T) \) is in both cases clearly "metal-like", i.e. decreasing with decreasing temperature, there are some peculiarities. First, the residual resistivity ratio (RRR) is high, with \( \rho(300)/\rho(0) \approx 2 \) for both materials (\( \rho(0) \) is the extrapolated zero temperature resistivity). Generally, the RRR is slightly higher for \( \text{Rb}_3\text{C}_6\text{O}_6 \) than for \( \text{K}_3\text{C}_6\text{O}_6 \), suggesting more "disorder" in the Rb-doped system. Second, \( \rho(T) \) shows unusual upward curvature over nearly the entire measured temperature range. Indeed, for both \( \text{K}_3\text{C}_6\text{O}_6 \) and \( \text{Rb}_3\text{C}_6\text{O}_6 \), \( \rho(T) \) is fit almost perfectly from just above \( T_c \) to well above room temperature by the functional form

\[
\rho(T) = A + BT^2. \tag{1}
\]

Is the "\( T^2 \)" fit meaningful? Very likely not. The large coefficient of thermal expansion in both materials implies that the density of states \( N(E_F) \) is highly temperature dependent\(^5\). An accurate analysis of \( \rho(T) \) thus necessitates a theoretical modeling of \( N(E_F) \), or, more directly, a measurement of \( \rho(T) \) under conditions of constant sample volume\(^4\), not constant pressure. Both approaches have been used\(^3,5\) and in each case \( \rho(T) \) no longer obeys Eq. (1). \( \rho(T) \) measured in \( \text{Rb}_3\text{C}_6\text{O}_6 \) under constant volume conditions is discussed in \$2.4\).

Although the functional forms of \( \rho(T) \) shown in Fig. 1 for \( \text{K}_3\text{C}_6\text{O}_6 \) and \( \text{Rb}_3\text{C}_6\text{O}_6 \) are quite reproducible from sample to sample, the absolute value of \( \rho(300K) \) is somewhat sample dependent. The "best" single crystal specimens (presumably those with the most homogeneous doping) have \( \rho(300K) \approx 1\text{m}\Omega\cdot\text{cm} \) and \( \rho(300K) \approx 2\text{m}\Omega\cdot\text{cm} \). Similar values are obtained for high quality films\(^6\). The uncertainties for the single crystal measurements are largely due to uncertainties in effective geometry.

Other methods have been applied to determine the absolute value of \( \rho \). Extrapolated infrared reflectance measurements by Rutter et al\(^7\) suggest for \( \text{K}_3\text{C}_6\text{O}_6 \) a low temperature resistivity \( \rho(0) = 0.4 \text{m}\Omega\cdot\text{cm} \), and microwave surface impedance measurements again on \( \text{K}_3\text{C}_6\text{O}_6 \) by Klein et al\(^8\) find \( \rho(0) = 0.5 \text{m}\Omega\cdot\text{cm} \). It is also possible to determine \( \rho(0) \) indirectly from an analysis of the fluctuation conductivity near \( T_c \) (see \$2.2); this yields\(^9\) \( \rho(0) = 0.12 \text{m}\Omega\cdot\text{cm} \) and \( \rho(0) = 0.22 \text{m}\Omega\cdot\text{cm} \).

Perhaps the best approach in determining \( \rho(0) \) is to apply the simple relation \( \rho(0) = 4\pi/\omega_c \• \chi(0) \). An analysis of upper critical field data yields\(^9,10\) values of the \( T=0 \) scattering time \( \tau \chi(0) = 1.7\pm0.5 \times 10^{-14} \text{s} \) and \( \tau \rho = 0.7\pm0.3 \times 10^{-14} \text{s} \). Combining this with theoretical values\(^11,12\) of the plasma frequency, \( 1.2\text{eV} \) and \( 1.1\text{eV} \) respectively for \( \text{K}_3\text{C}_6\text{O}_6 \) and \( \text{Rb}_3\text{C}_6\text{O}_6 \), we find \( \rho(0) = 0.18\pm0.06 \text{m}\Omega\cdot\text{cm} \) and \( \rho(0) = 0.57\pm0.21 \text{m}\Omega\cdot\text{cm} \).

These values lie intermediate to those determined from direct single crystal and film dc resistivity measurements and those extracted from a fluctuation conductivity analysis; they imply "ideal sample" values for the room temperature resistivity \( \rho(300K) \approx 0.4 \text{m}\Omega\cdot\text{cm} \) and \( \rho(300K) \approx 1.1 \text{m}\Omega\cdot\text{cm} \).

### 2.2 Low Temperature dc Resistivity: Paraconductivity Near \( T_c \)

In the normal state just above the thermodynamic superconducting phase transition temperature \( T_c \), order parameter fluctuations can affect the normal state resistivity \( \rho(T) \). In a conventional isotropic superconductor, such fluctuations are normally not observed since the temperature interval over which they exist is too small. The large intrinsic disorder in \( \text{K}_3\text{C}_6\text{O}_6 \) and \( \text{Rb}_3\text{C}_6\text{O}_6 \) extends the fluctuation temperature range such that it is experimentally accessible. In fact, \( \text{A}_3\text{C}_6\text{O}_6 \) superconductors have provided the first direct experimental test\(^2\) of 3-D superconductivity fluctuation theory\(^13\), dating to 1968!

Fig. 2 shows the dc resistivity of \( \text{Rb}_3\text{C}_6\text{O}_6 \) in the immediate temperature range above \( T_c = 30.2\text{K} \). The solid line is an extrapolation of the normal state \( \rho(T) \); over this small temperature range a linear fit is adequate. The deviation of the experimental data from the extrapolated curve reflects the excess fluctuation conductivity \( \sigma' \) (i.e. paraconductivity), which from Aslamazov-Larkin theory\(^13\) may be expressed as\(^14\)

\[
\sigma = \sigma_{\text{exc}} t^{-(4-D)/2} \tag{2}
\]

where \( \sigma_{\text{exc}} = 2\pi e^2/32h\xi(0) \), \( t = (T-T_c)/T_c \), \( D \) is the dimensionality of the system, and \( \xi(0) \) is the superconducting coherence length. The inset to Fig. 2 shows on a log-log plot the experimentally determined fluctuation conductivity \( \sigma' \) versus the reduced temperature \( t \). A good fit is obtained for \( D=3 \), which indicates \( \text{Rb}_3\text{C}_6\text{O}_6 \) is an intrinsic three dimensional superconductor. Similar results are obtained for \( \text{K}_3\text{C}_6\text{O}_6 \), where again \( D=3 \) provides the best fits. However, in \( \text{K}_3\text{C}_6\text{O}_6 \) there is evidence for an additional Maki-Thompson term\(^15\) in the fluctuation conductivity with a pair-breaking parameter \( \delta \approx 0.58 \). In \( \text{Rb}_3\text{C}_6\text{O}_6 \), the Maki-Thompson term can be neglected implying a larger \( \delta \), which in turn suggests a larger electron-phonon coupling in \( \text{Rb}_3\text{C}_6\text{O}_6 \).

The fluctuation conductivity can be used to evaluate the intrinsic normal state resistivity, in a manner that eliminates sample
Is the characteristic scattering length scale in $A_2C_60$ the crystal lattice constant ($-14\,\text{Å}$) or the on-ball carbon-carbon distance ($-1\,\text{Å}$)?

Fig. 3 shows $\rho(T)$ measured up to 800K in single crystal $K_3C_60$ and up to 700K in single crystal $Rb_3C_60$. A pulsed heating technique has been used to prevent de-intercalation of the alkali metals at high temperature. $\rho(T)$ in $K_3C_60$ accurately follows a $T^2$ temperature dependence all the way to 800K, while the data for $Rb_3C_60$ show distinct curvature changes at high temperatures. Do these data sets indicate resistivity saturation? Naively one might expect resistivity saturation to manifest itself as a high temperature plateauing of $\rho(T)$; such behavior is clearly not observed for either material. However, the special properties of the alkali-fullerides necessitate a more careful theoretical treatment to confirm or exclude the observation of resistivity saturation.

As the electronic mean free path approaches the characteristic lattice spacing, Bloch-Boltzmann transport theory fails. One accounts for this effect by imposing a phenomenological minimum electron scattering time $\tau_{\text{sat}}$, which corresponds to a length scale $\ell_{\text{sat}} = \frac{\tau_{\text{sat}}}{v_F}$, on the order of the interatomic spacing. This minimal time acts as an offset to the Poisson distribution of electron scattering events, yielding a parallel resistor form $^{17,18}$ for the measured resistivity

$$\rho^{-1} = \rho_B^{-1} + \rho_{\text{sat}}^{-1}$$

where $\rho_B$ is the Bloch-Boltzmann resistivity; it is a sum of the residual resistivity $\rho_0$ and the temperature-dependent electron-phonon resistivity $\rho_{\text{ep}}$. $\rho_{\text{ep}}$ given by the Ziman resistivity formula $^{19}$, incorporates the phonon spectrum and electron-phonon coupling. Furthermore, $\rho_{\text{ep}}$ is proportional to $|N(E_F)|^2$. The strong temperature dependence of the density of states in $A_2C_60$ conductors must be taken into account in an analysis of the resistivity.

Hou et al $^{16}$ have performed a careful analysis of the high temperature resistivity of $K_3C_60$ and $Rb_3C_60$, by applying Eq. (3). A temperature-dependent density of states is used and different model forms of the phonon spectral function are tested. Fig. 4 compares the results of the theoretical analysis to the experimental data. For $Rb_3C_60$, a good fit is obtained using a model coupling spectrum $^{19}$ with average phonon frequency $<\omega>$ = 2000K (a fit using a model coupling spectrum $^{20}$ with $<\omega>$ = 500K yields an inferior fit). The analysis indicates resistivity saturation and yields $\rho_{\text{sat}} = 6.3\pm2.4\,\text{mΩ}\cdot\text{cm}$. With $v_F = 1.6\times10^7\,\text{cm/sec}$, we obtain $\ell_{\text{sat}} = 1.0\pm0.5\,\text{Å}$. In $Rb_3C_60$, the saturation length is on the order

2.3 High Temperature dc Resistivity: Saturation

For a metal at high temperature, one might expect the resistivity to saturate as the scattering length (i.e. electronic mean free path) approaches the intrinsic interatomic distance. The high temperature resistivity of $K_3C_60$ and $Rb_3C_60$ has been examined $^{6,16}$ with this in mind.
Figure 3: Resistivity of $K_3C_{60}$ and $Rb_3C_{60}$ over an extended temperature range. Different symbols correspond to different pulsed heating runs. Data below 300K (dashed lines) were obtained using conventional slow cooling methods.

Figure 4: Resistivity versus temperature for $Rb_3C_{60}$ (upper curves) and $K_3C_{60}$ (lower curves). Open circles represent experimental data normalized to $\rho(0) = 0.18$ m$\Omega$-cm and $\rho(Rb) = 0.57$ m$\Omega$-cm. Shaded regions indicate uncertainties in absolute resistivity normalization. The solid and dashed lines are theoretical fits$^{16}$ from which resistivity saturation parameters are extracted. For $Rb_3C_{60}$ the dashed line fit uses the electron-plasmon coupling spectrum from Jishi et al.$^{20}$ ($<\omega>$ = 500K), while the solid line fit uses the spectrum of Varma et al.$^{19}$ ($<\omega>$ = 2000K). Both models include a temperature dependent density of states. The fits to the $K_3C_{60}$ data use the coupling spectrum from Schluter et al.$^{21}$ ($<\omega>$ = 1200K). The solid line includes a temperature-dependent density of states whereas the dashed fit does not.
of the carbon-carbon bond length. Fig. 4 also shows the theoretical fit to the K$_3$C$_60$ data. The data do not show obvious signs of resistivity saturation, but the analysis, using a model coupling spectrum with $\omega$ = 1200K, yields a lower bound on the magnitude of the saturation resistivity: $\rho_{\text{sat}} > 3\text{m}\Omega\cdot\text{cm}$ and $\xi_{\text{sat}} < 1.5\AA$. The characteristic saturation length in K$_3$C$_60$ is expected to be again on the order of the carbon-carbon bond length.

The high-temperature resistivity measurements and analysis suggest that the single particle states relevant to transport at high temperatures in A$_3$C$_60$ have a length scale much smaller than the fcc lattice constant or the C$_60$ ball nearest neighbor center-to-center distance. Although the extended-state description of electronic dynamics may be applicable to A$_3$C$_60$ at low temperatures, the small saturation scattering length implies that the character of the appropriate electronic states at high temperatures is quite different.

2.4 Pressure Studies: Constant-volume $\rho(T)$

As mentioned above, the large thermal expansion of A$_3$C$_60$ complicates direct comparison of transport measurements with theoretical model predictions. Model predictions usually assume a constant sample volume, while most experiments are carried out under constant pressure conditions. Crespi et al. and Hou et al. have compared experimental constant-pressure $\rho(T)$ data of K$_3$C$_60$ and Rb$_3$C$_60$ to various electron-phonon scattering models incorporating a temperature-dependent density of states $N(E_F)$. One such comparison has been shown in Fig. 4.

Here we adopt a different strategy: a direct experimental determination of the constant-volume $\rho(T)$, which we designate $\rho_v(T)$. This is possible by applying external, hydrostatic pressure to the sample to prevent thermal expansion. In an ideal experiment, the pressure would be adjusted at each temperature so as to compensate for thermal expansion and bring the sample volume to a chosen reference volume $v_0$, say the ambient pressure volume at 4.2K. In practice the voltage of the sample is not measured directly but is inferred from the applied pressure and the independently determined (temperature dependent) bulk modulus. In addition, in a typical experiment $\rho(T)$ is determined as a function of temperature only for a series of semi-temperature-independent "clamp" pressures; the data are then interpolated to determine $\rho_v(T)$.

Fig. 5 shows the result of $\rho_v(T)$ determined for a Rb$_3$C$_60$ crystal. The maximum pressure applied, at 350K, is about 8kbar. From 100K to 350K, the $\rho_v(T)$ data differ significantly from the "constant pressure" $T^2$ result of Eq. (2). Instead, for constant sample volume a linear temperature dependence is observed, i.e.

$$\rho_v(T) = a + bT. \quad (4)$$

This is the "intrinsic" resistivity of Rb$_3$C$_60$. The inset to Fig. 5 shows $\rho_v(T)$ over an extended temperature range. The data near $T_c$ are not reliable due to uncertainties in the data calibration at high temperatures.

In the higher temperature regime, the slope of the $\rho_v(T)$ curve, $dp/dT$, is directly related to the transport electron-phonon coupling constant $\lambda_{tr}$. Combining $\rho(T) = 4\pi/\omega_p^2\tau$ with the first term in the scattering time expansion (valid for $T<80K$), $\tau^{-1} = \lambda_{tr} k_B T / h$, we obtain

$$\lambda_{tr} = (\hbar \omega_p^2 / 16 \pi^3 k_B) dp/dT = 0.246 (\hbar \omega_p / 2\pi^2) dp/dT. \quad (5)$$
the later equality holding with units \[ \frac{dp}{dT} = \mu \Omega \cdot \text{cm}/K \text{ and } [\hbar \omega_p/2 \pi] = \text{eV}. \] The relative slope of \( \rho_\nu(T) \) in Fig. 5 is \((1/p_0)\frac{dp}{dT} = (1.9 - 2.2) \times 10^{-3} \text{ K}^{-1}. \) To obtain \( \frac{dp}{dT} \), \( \rho_0, \) i.e. \( \rho_{RB}(295K) \), must be known. From §2.1 we have \( \rho_0 = 1.1 \mu \Omega \cdot \text{cm} \), which, with a plasma frequency of 1.2eV, leads to \( \lambda_T = 0.30 - 0.35. \) The directly measured crystal \( \rho_{RB}(295K) = 2.5 \mu \Omega \cdot \text{cm} \) leads to \( \lambda_T = 0.65 - 0.80. \)

It is interesting to note that, identifying \( \lambda_T \) with the superconductivity electron-phonon coupling constant \( \lambda \), the evaluated range of \( \lambda_T \)'s is entirely consistent with a 30K superconducting transition temperature, assuming a conventional phonon-mediated electron-electron pairing mechanism with characteristic phonon frequencies of order 2000K. It is not necessary to invoke "exotic" scattering or superconductivity mechanisms to account for the measured \( \rho(T) \) or \( T_c \) in \( \text{Rb}_3\text{C}_{60}. \)

2.5 Hall Effect

For \( \text{A}_3\text{C}_{60} \), we might expect that the intercalated alkali atoms donate one electron each to the host \( \text{C}_{60} \) structure, leading to 3 free electrons per \( \text{C}_{60}. \) In the most elementary (free electron) model, the Hall coefficient is given by \( R_H = -1/n, \) with \( n \) the carrier concentration. With 3 electrons/\( \text{C}_{60}, \) \( R_H = -1.4 \times 10^{-9} \text{ m}^3/\text{C}. \)

Fig. 6 shows \( R_H \) \( (T) \) determined for single crystal \( \text{K}_3\text{C}_{60} \) and \( \text{Rb}_3\text{C}_{60}. \)

For \( \text{K}_3\text{C}_{60}, \) \( R_H \) at room temperature is \( -1.4 \times 10^{-9} \text{ m}^3/\text{C}, \) which coincides with the free electron result with 3 electrons/\( \text{C}_{60}. \) This coincidence is most likely fortuitous; the free electron model is of limited utility in a system with a complex Fermi surface. Over the measured temperature range, \( R_H \) for \( \text{K}_3\text{C}_{60} \) is negative and increases linearly with increasing temperature with a slope \( 0.0055 \times 10^{-9} \text{ m}^3/\text{CK}. \) For \( \text{Rb}_3\text{C}_{60}, \) \( R_H \) again increases linearly with temperature with the same slope; the data are displaced upward from those of \( \text{K}_3\text{C}_{60} \) and there is a zero crossing at \( T = 120K. \)

The observation of a strong influence of thermal expansion on the dc electrical resistivity of \( \text{A}_3\text{C}_{60} \) suggests that lattice thermal expansion should be taken into account in a discussion of \( R_H(T). \) The lattice thermal expansion of \( \text{A}_3\text{C}_{60} \) in the temperature range of interest can be adequately expressed as \( a(T) = a_0 + bT, \) where \( b = 3.52 \times 10^{-4} \text{ Å/K for } \text{K}_3\text{C}_{60} \) and \( 4.4 \times 10^{-4} \text{ Å/K for } \text{Rb}_3\text{C}_{60}. \)

Fig. 7 shows \( R_H \) versus lattice constant \( a. \) This striking relationship

![Figure 6: Hall coefficient for \( \text{K}_3\text{C}_{60} \) and \( \text{Rb}_3\text{C}_{60}. \)](image)

![Figure 7: Hall coefficient for \( \text{K}_3\text{C}_{60} \) and \( \text{Rb}_3\text{C}_{60} \) plotted as a function of lattice constant. A universal function for \( R_H \) is observed, suggesting that \( R_H \) at constant sample volume, is temperature and material independent.](image)

shows that the difference in \( R_H \) between \( \text{K}_3\text{C}_{60} \) and \( \text{Rb}_3\text{C}_{60} \) at a given temperature can be ascribed purely to the difference in the lattice constant between the two materials, and that the temperature dependence of \( R_H \) in the two materials can also be ascribed purely to changes in the lattice constant. In other words, \( R_H \) at constant sample volume is temperature
independent for both materials. For both $K_3C_{60}$ and $Rb_3C_{60}$, the experimentally determined Hall coefficient can be expressed as

$$R_H = [-2.8 + 13.9(a - 14.15)] \times 10^{-9} \text{ m}^3/\text{C},$$

where $a$ is measured in angstroms. Eq. (6) is expected to hold for all $A_3C_{60}$ conductors, including mixed alkali alloys, although this prediction has not been tested.

Prior calculations by Erwin and Pickett of the Hall coefficient in $A_3C_{60}$ using the Jones-Zener solution of the Boltzmann equation in an orientationally ordered system at $T=0$ with isotropic scattering time yield $R_H = 7 \times 10^{-9} \text{ m}^3/\text{C}$ for both $K_3C_{60}$ and $Rb_3C_{60}$. This result, which is weakly pressure dependent, arises from a complex weighted average of curvature over the Fermi surface. However, this prediction is inconsistent with the data of Figs. 6 and 7. The experiments show $R_H$ to be smaller in absolute magnitude and negative over a wide temperature range. Furthermore, Eq. (6) suggests a pressure dependence to $R_H$ much stronger than that calculated. Since $R_H$ is sensitive to details of the Fermi surface curvature, it is likely that the discrepancy between theory and experiment is due to disorder-induced modification of the Fermi surface (not taken into account in the orientationally ordered calculation).

Lu et al. have modeled the effects of lattice-constant-dependent disorder broadening at the Fermi level. The variation in $R_H$ with lattice constant is accounted for by assuming that the relative k-space extent of the disorder broadening of the Fermi surface increases with increasing lattice constant. If the effects of the disorder have a characteristic energy scale which is to lowest order independent of the lattice constant, then the k-space extent of the disorder broadening will be an increasing function of the density of states at the Fermi level and hence an increasing function of the lattice constant. This goes hand in hand with the larger dc residual resistivity of $Rb_3C_{60}$ compared to $K_3C_{60}$ (§2.1). The behavior is similar to that suggested for disordered nearly free electron metals, wherein $R_H$ is inversely dependent on the density of states.

The model of Lu et al predicts that the Hall coefficient $R_H$ in $A_3C_{60}$ will increase from negative values at low disorder broadening to positive values at larger disorder broadening. This is consistent with the experimental results, which show negative $R_H$ for the $K_3C_{60}$ sample at low $T$ (small density of states) going over to positive $R_H$ for the $Rb_3C_{60}$ sample at high $T$ (large density of states). A simplified calculation applying the disorder-broadening model predicts the correct ratio of the extrapolated $T=0$ Hall coefficients of $K_3C_{60}$ and $Rb_3C_{60}$.

2.6 Thermoelectric Power

For a conventional metal, the thermoelectric power (TEP) is expected to be small in magnitude (order $\mu\text{V/K}$), have a sign reflecting the sign of the charge carriers, and be linear in temperature.

Fig. 8 shows the TEP for $K_3C_{60}$ and $Rb_3C_{60}$ as measured by Inabe et al. These crystals have been grown from a $CS_2$ solution and they display an anomalously high resistivity, suggesting incomplete doping or $CS_2$ impurities. However, the TEP results are consistent with those obtained on vapor-transport grown crystals. The data of Fig. 8 indicate a linear TEP for both materials, with room temperature values $-10 \mu\text{V/K}$ for $K_3C_{60}$ and $-15 \mu\text{V/K}$ for $Rb_3C_{60}$. The hump structure near 50K in both materials may be due to phonon drag. The TEP results are consistent with metallic behavior.

![Figure 8: Thermoelectric power versus temperature for $K_3C_{60}$ and $Rb_3C_{60}$.](image-url)
Inabe et al have used the TEP results to estimate the density of states ratio between Rb$_2$C$_{60}$ and K$_3$C$_{60}$; they obtain $\frac{N_{Rb}(Ep)}{N_K(Ep)} = 1.5-1.8$. Band structure calculations indicate that the density of states ratio is $\sim 1.2$.

### 3 Transport in A$_1$C$_{60}$

The recently discovered polymerized A$_1$C$_{60}$ materials have been characterized by a variety of techniques, including structural, Raman, IR, and spin resonance. The structure is intriguing in that it consists of C$_{60}$ molecules covalently bonded into linear chains. Although such a geometry is a template for quasi-one-dimensional conduction, the actual charge transport in these materials may be quite complicated. For example, the C$_{60}$-C$_{60}$ bond is expected to be a serious impediment to charge conduction, and transport along the chain axis may actually follow a "zig-zag" path where charges hop from one polymerized chain to another.

Clearly transport measurements in single crystal specimens would be highly desirable. Unfortunately, even the best available samples of A$_1$C$_{60}$ (A=K,Rb,Cs) are not true single crystals, so that the polymerization direction is not uniform throughout the macroscopic specimen. Hence all transport measurements to date on the A$_1$C$_{60}$ materials represent an average over different crystal directions and cannot distinguish crystal anisotropies. Nevertheless, such measurements are illuminating and show a rich electronic structure.

#### 3.1 Dc Resistivity of A$_1$C$_{60}$

Figure 9 shows the dc electrical resistivity $\rho(T)$ of K$_1$C$_{60}$ from room temperature to 4K. Between 300K and 50K a metal-like behavior is observed, with $\rho(T)$ decreasing smoothly with decreasing temperature. However, $\rho(T)$ does not follow any conventional form (the solid line in the figure is fit to the functional form $A + BT + CT^2$). It is possible that thermal expansion plays a role in the observed $\rho(T)$, although in this covalently-bonded polymer the effect is expected to be less severe than in A$_3$C$_{60}$ compounds.

Below 50K, the resistivity of K$_1$C$_{60}$ is anomalous and increases with decreasing temperature. This may be the signature of a structural and/or electronic phase transition. For example, the spin density wave (SDW) transition in the low dimensional organic charge transfer salt (TMTSF)$_2$PF$_6$ manifests itself as a very similar upturn in $\rho(T)$.

The low temperature resistivity in K$_1$C$_{60}$ does not, however, follow any activated form as might be expected with the opening of a gap at the Fermi energy. If the 50K transition in A$_1$C$_{60}$ is associated with formation of a SDW (or, alternatively, a charge density wave), then the Fermi surface is most likely only partially gapped. In K$_1$C$_{60}$, no superconductivity is observed down to 1.9K under ambient pressure conditions. In addition, no magnetoresistance is observed (using fields up to 7T) at any temperature.

Figure 10 shows the dc resistivity for Rb$_1$C$_{60}$ and Cs$_1$C$_{60}$. For both materials, the transport properties indicate semiconducting, rather than metallic, behavior. Activated behavior with a temperature-independent activation energy is not observed. A careful examination of the $\rho(T)$ curve
3.2 Pressure Effects on the dc Resistivity

The dc resistivity of K$_2$C$_{60}$ is extremely pressure dependent. At room temperature, an applied hydrostatic pressure of 15kbar reduces the resistivity by 70%$^{24}$. A further increase in pressure has little effect. This suggests a saturation in the charge transfer matrix element (possibly inter-chain hopping). The electrical resistivity measured$^{24}$ under such “saturation” conditions is shown in Fig. 11. A fully conventional metallic $\rho(T)$ is observed, with no evidence for a low temperature phase transition. No superconductivity is observed to 4.2K (14.5kbar). It is likely that applied pressure makes K$_2$C$_{60}$ more three-dimensional, hence suppressing the quasi 1-D low temperature magnetic phase transition.

Pressure studies have also been recently been performed$^{36,37}$ on Rb$_2$C$_{60}$ and Cs$_2$C$_{60}$. The semiconducting $\rho(T)$ is transformed into a metallic $\rho(T)$, with sharp "semiconductor-to-metal" transitions observed with decreasing temperature at fixed pressures. These transitions suggest a rich electronic/magnetic phase diagram for the A$_2$C$_{60}$ polymers. Additional transport and structural measurements, if possible on single domain samples, would be highly desirable.

Figure 11: High-pressure dc electrical resistivity of K$_2$C$_{60}$. Metallic behavior is observed over the entire measured temperature range.
Acknowledgments

This work was performed in collaboration with G. Briceno, M.L. Cohen, V. H. Crespi, M.F. Fuhrer, J. Hone, J.G. Hou, K. Khazeni, L. Lu, W. Vareka, and X.-D. Xiang. Support is acknowledged from NSF Grant DMR-94-04755 and DOE grant DE-AC03-76F00098. AZ received support from the Miller Institute for Basic Research in Science.

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