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Q. Jiang

Z.H. Zhang

X.B. Wang

T. Chui

and

D.C. Tian

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Q. Jiang[1]
International Centre for Theoretical Physics, Trieste, Italy,

Z.H. Zhang, X.B. Wang
Department of Physics, Wuhan University, Wuhan 430072, People's Republic of China,

T. Chui
Bartol Research Institute, University of Delaware, Newark DE19716, USA

and

D.C. Tian
Department of Physics, Wuhan University, Wuhan 430072, People's Republic of China,
International Center for Materials Physics, Shenyang 110015, People's Republic of China
and
Bartol Research Institute, University of Delaware, Newark DE19716, USA.

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[1]Permanent address: Department of Physics, Wuhan University, Wuhan 430072, People's Republic of China.

ABSTRACT

Based on the electron–fracton interaction, the Hall effect and resistivity in the normal states of high-T_c oxides are calculated by solving the quantum transport equations. The numerical results show that, under certain conditions, the inplane resistivity exhibits a nearly linear temperature dependence and the Hall angle exhibits a quadratic temperature dependence.

Like the unusual nearly linear temperature dependence of the in-plane resistivity ρ_{xx} in the "normal" metallic state of the cuprate superconductors , the strong temperature dependence of the Hall effect is also puzzling anomalous. Experimental measurements find that the reciprocal of the Hall coefficient, $1/R_H$ is approximately linear in temperature and furthermore the Hall angle $cot\theta_H$ varies as T^2 for many cuprate systems [1]. The Hall angle $cot\theta_H$ persists this power-law behavior even up to 500K, and even in samples with higher impurity concentrations and reduced carrier densities for which R_H and ρ_{xx} do not have simple T dependencies [2].

The measurements demonstrate an apparent inconsistency that lies in the explanation of the Hall data based on a usual Drude single model with isotropic scattering and the conventional Boltzmann transport equations. Strong evidence exists also against other scenarios based on multiple Drude bands or magnetic skew-scattering mechanism [3]. Much effort has been devoted to the explanation of the puzzling and striking anomalies such as a transverse relaxation rate due to spinon-spinon scattering [4], the Fermi-liquid picture combined with an anisotropic scattering [5], cyclotron resonance in a Drude model [6], a variant of bipolaron model [7] and localized fermions [8] and so on.

Shortly after the discovery of high-T_c cuprates superconductors, Butter and Blumen first suggested a possible explanation for the high-T_c superconducting phase in the Y-Ba-Cu-O system in the electron-fracton model [9]. Assuming that the distribution of oxygen deficiencies, which exist commonly in most high-T_c cuprates, is strongly disordered in the copper oxide planes, we can construct a two-dimensional percolation network to simulate the distribution of oxygen holes. The doping oxygen holes may affect significantly the distribution of the force constants related to the atomic vibration. As we know, percolation networks appear to be homogeneous at length scales L longer than the percolation correlation length ξ, and thus support propagating phonons. For shorter length scales $L < \xi$, however, the random networks exhibit

fractal characteristic so that one would expect localized vibrational excitations called fractons [10]. The vibrational excitations on such a structure consist of phonons below a crossover frequency ω_c and fractons over ω_c. In our previous works, we have studied the temperature-dependence of the resistivity arising from the scattering of conduction electrons off fractons and the results showed that this contribution to the resistivity is nearly linear in temperature [11]. We also present a quantitative calculation of the transition temperature T_c and gap parameter of superconductors with fractal structure [12].

The purpose of this paper is to investigate the contribution of electron-fracton interaction to the in-plane conductivity σ_{xx} and to the Hall conductivity σ_{xy} (magnetic field $\vec{H}\|\vec{c}$) by solving quantum transport equations proposed by Mahan and co-worker [13]. We expect to see whether in-plane resistivity ρ_{xx} varies as T and σ_{xy} varies as T^{-3} under certain conditions, and to find whether the interaction mechanism may offer a possible explanation to the anomalous Hall effect in the normal state of high -T_c cuprates.

For steady-state transport phenomena in a metal, a many-body quantum transport equation including a magnetic field derived by Hansch and Mahan[13] is

$$e\vec{E}\cdot\vec{v}_q Im\Sigma^r(\vec{q},\omega)\frac{\partial n_F(\omega)}{\partial\omega}A^2(\vec{q},\omega)+\frac{ie}{c}\vec{v}_q\cdot\vec{H}\times\nabla_{\vec{q}}G^< = \Sigma^> G^< - \Sigma^< G^> \quad (1)$$

where $n_F(\omega)$ is an equilibrium Fermi distribution and $\vec{v}_q$ is a group velocity of electrons given by $\nabla_{\vec{q}}\epsilon_{\vec{q}}$, $G^{<>}$ are the Kadanoff-Baym physical Green-functions. $\Sigma^>$ reprents the rate of conduction electrons scattered to leave from the configuration of energy ω and quasi-momentum $\vec{q}$, $\Sigma^<$ is regarded as the rate of conduction electrons scattered to arrive the corresponding energy-momentum configuration. In the case of one-fracton processes, they can be expressed as

$$\Sigma^<(\vec{q},\omega) = \sum_{\vec{k}} |g_{\vec{k}}|^2[(n_B(\omega)+1)G^<(\vec{q}+\vec{k},\omega+\omega_{\vec{k}}) + n_B(\omega)G^<(\vec{q}+\vec{k},\omega-\omega_{\vec{k}})]$$

$$\Sigma^>(\vec{q},\omega) = \sum_{\vec{k}} |g_{\vec{k}}|^2[n_B(\omega)G^>(\vec{q}+\vec{k},\omega+\omega_{\vec{k}}) + (n_B(\omega)+1)G^>(\vec{q}+\vec{k},\omega-\omega_{\vec{k}})] \quad (2)$$

where $g_{\vec{k}}$ is electron-fracton scattering matrix element given by[11]

$$|g_{\vec{k}}|^2 = \frac{\hbar\lambda_0^2}{2B}\frac{|\vec{k}|^{2-D}}{\omega_\lambda^{\vec{k}}} \quad (3)$$

λ_0 is the electron-fracton coupling constant and $\omega_\lambda^{\vec{k}}$ is fracton frequency, D is the fractal dimensionality.

The Eq.(1) is obtained by assuming that the metal has a symmetrical Fermi surface and we can make a standard approximation that the retarded self-energies $\Sigma^r(\vec{q},\omega)$ have a negligible dependence upon wave vector $\vec{q}$ and can use the shorthand notation $\Gamma(\omega) = -Im\Sigma^r(\vec{q},\omega)$. Solving Eq.(1) self- consistently with an ansatz:

$$G^<(\vec{q},\omega) = iA(\vec{q},\omega)n_F(\omega) - i\frac{\partial n_F(\omega)}{\partial\omega}\gamma(\vec{q},\omega)$$

$$G^>(\vec{q},\omega) = G^<(\vec{q},\omega) - iA(\vec{q},\omega) \quad (4)$$

we may deduce a simplified version

$$e\vec{E}\cdot\vec{v}_q\Gamma(\omega)A^2(\vec{q},\omega) + \frac{e}{c}\vec{H}\times\vec{v}_q\cdot\nabla_{\vec{q}}\gamma(\vec{q},\omega) = 2\Gamma(\omega)\gamma(\vec{q},\omega)$$

$$-A(\vec{q},\omega)\sum_{\vec{k}} g_{\vec{k}}^2[N_-\gamma(-) + N_+\gamma(+)] \quad (5)$$

where $\gamma(\vec{q},\omega)$ is unknown vertex function and the other notations in Eq.(5) are

$$N_- = n_B(\omega_{\vec{k}}) + 1 - n_F(\omega-\omega_{\vec{k}})$$

$$N_+ = n_B(\omega_{\vec{k}}) + n_F(\omega+\omega_{\vec{k}})$$

$$\gamma(-) = \gamma(\vec{q}+\vec{k},\omega-\omega_k)$$

$$\gamma(+) = \gamma(\vec{q}+\vec{k},\omega+\omega_k)$$

When the width of the Landau levels is much less than the Fermi energy, a quasi-classical treatment of the effect of the higher reduced field is sufficient. Then the group velocity of electrons $\vec{v}_q$ in a weak field must satisfy $\vec{v}_q = \vec{q}/m^*$, m^* is the effective mass of electron. Thus following the similar treatment of Yi [14], we can expand the vertex function $\gamma(\vec{q},\omega)$ in the order of the dimensionless parameter $\vec{h}(\omega) = \frac{e\vec{H}}{2m^*\Gamma(\omega)c}$

$$\gamma(\vec{q},\omega) = \sum_{n=0}^{\infty} h^n(\omega)\gamma^{(n)}(\vec{q},\omega) \quad (6)$$

Instituting Eq.(6) to Eq.(5) and comparing the order of $h(\omega)$ in the two sides of the equation, we have

$$\gamma^{(0)}(\vec{q},\omega) = \frac{e\vec{E}\cdot\vec{q}A^2(\vec{q},\omega)}{2m^*} + \frac{A(\vec{q},\omega)}{2\Gamma(\omega)}\sum_{\vec{k}} g_{\vec{k}}^2[N_-\gamma^{(0)}(-) + N_+\gamma^{(0)}(+)] \quad (7)$$

and

$$\gamma^{(n)}(\vec{q},\omega) = \vec{e}_h\times\vec{q}\cdot\nabla_{\vec{q}}\gamma^{(n-1)}(\vec{q},\omega) + \frac{A(\vec{q},\omega)}{2\Gamma(\omega)}\sum_{\vec{k}} g_{\vec{k}}^2[N_-\gamma^{(n)}(-) + N_+\gamma^{(n)}(+)] \quad (8)$$

for n=1, 2, 3

where $\vec{e}_h$ is the unit vector of magnetic field $\vec{H}$. $\gamma^{(0)}(\vec{q},\omega)$ in Eq.(7) is corresponding to the zero-magnetic-field vertex function and it has been solved by Hansch in the Ref [15]

$$\gamma^{(0)}(\vec{q},\omega) = \frac{e}{m^*}A(\vec{q},\omega)\Lambda(\vec{q},\omega)\vec{E}\cdot\vec{q} \qquad (9)$$

where $\Lambda(\vec{q},\omega) = \frac{1}{2}(A(\vec{q},\omega)+\frac{\Lambda(\omega)-1}{\Gamma(\omega)})$ is another unknown function and $\Lambda(\omega)$ is a function determined by the integral equation

$$\Lambda(\omega) = 1 + \pi\int_{\omega_c}^{\omega_{FD}} du[(\alpha^2F(u) - \alpha_{tr}^2F(u))\frac{\Lambda(u+\omega)}{\Gamma(u+\omega)}(n_B(u) + n_F(u+\omega))$$

$$+ \frac{\Lambda(u-\omega)}{\Gamma(u-\omega)}(n_B(u) + n_F(u-\omega))] \qquad (10)$$

The functions $\alpha^2F(u)$ and $\alpha_{tr}^2F(u)$ under the Debye approximation have been calculated in the previous paper [11]

$$\alpha^2F(u) = Cu^{d/D-2}$$

$$\alpha_{tr}^2F(u) = C_{tr}u^{3\tilde{d}/D-2} \qquad (11)$$

where $\tilde{d}$ is the fracton dimensionality.

The next step is to solve the first-order term of vertex function, $\gamma^{(1)}(\vec{q},\omega)$. Inserting Eq.(9) into Eq.(8) and setting n=1, we have

$$\gamma^{(1)}(\vec{q},\omega) = \vec{e}_h\times\vec{q}\cdot\nabla_{\vec{q}}(\frac{e}{m^*}A\Lambda\vec{E}\cdot\vec{q}) + \frac{A}{2\Gamma(\omega)}\sum_{\vec{k}}g_{\vec{k}}^2[N_-\gamma^{(1)}(-) + N_+\gamma^{(1)}(+)] \qquad (12)$$

For the lowest-order approximation, $(\vec{e}_h\times\vec{q})\cdot\nabla_{\vec{q}}(\frac{e}{m^*}A\Lambda\vec{E}\cdot\vec{q}) = \frac{e}{m^*}A\Lambda\vec{E}\cdot(\vec{e}_h\times\vec{q})$

Generally, for a unknown function $X(\vec{q},\omega)$ which satisfy the following intergal equation

$$X(\vec{q},\omega) = F(\vec{q},\omega) + \frac{A}{2\Gamma(\omega)}\sum_{\vec{k}}g_{\vec{k}}^2[N_-X(-) + N_+X(+)] \qquad (13)$$

where $F(\vec{q},\omega)$ is given, comparing the Eq.(12) with the Eq.(7) and its solution Eq.(9), we can write the formal solution of $X(\vec{q},\omega)$ as

$$X(\vec{q},\omega) = F(\vec{q},\omega)\frac{\Lambda(\vec{q},\omega)}{A(\vec{q},\omega)} \qquad (14)$$

Therefore, we obtain

$$\gamma^{(1)}(\vec{q},\omega) = \frac{e}{m^*}A(\vec{q},\omega)\Lambda(\vec{q},\omega)\vec{E}\cdot(\vec{e}_h\times\vec{q}) \qquad (15)$$

and similarly, the solution of the vertex function of arbitrary order is

$$\gamma^{(2n)}(\vec{q},\omega) = (-1)^n\gamma^{(0)}(\vec{q},\omega)$$

$$\gamma^{(2n+1)}(\vec{q},\omega) = (-1)^n\gamma^{(1)}(\vec{q},\omega) \qquad (16)$$

for n=1,2,3... ,

Then, the final solution of the vertex function is

$$\gamma(\vec{q},\omega) = \sum_{n=0}^{\infty}h^n(\omega)\gamma^{(n)}(\vec{q},\omega) = \frac{1}{1+h^2(\omega)}\gamma^{(0)}(\vec{q},\omega) + \frac{h(\omega)}{1+h^2(\omega)}\gamma^{(1)}(\vec{q},\omega) \qquad (17)$$

Knowing the vertex function $\gamma(\vec{q},\omega)$ which is linear in the electric field, we can calculate the current density

$$\vec{j} = e^2\int\frac{d^3\vec{q}}{(2\pi)^3}\int\frac{d\omega}{2\pi}(-\frac{\partial n_F(\omega)}{\partial\omega})\vec{v}_{\vec{q}}\gamma(\vec{q},\omega) \qquad (18)$$

The longitudinal σ_{xx} and the transverse (Hall) σ_{xy} conductivities are then

$$\sigma_{xx} = \sigma_{yy} = \frac{2e^2n_0}{m^*}\int\frac{d\omega}{2\pi}(-\frac{\partial n_F}{\partial\omega})\frac{1}{1+h^2(\omega)}\int\frac{d^3\vec{q}}{(2\pi)^3}q_x^2A(\vec{q},\omega)\Lambda(\vec{q},\omega)$$

$$\sigma_{xy} = -\sigma_{yx} = \frac{2e^2n_0}{m^*}\int\frac{d\omega}{2\pi}(-\frac{\partial n_F}{\partial\omega})\frac{h(\omega)}{1+h^2(\omega)}\int\frac{d^3\vec{q}}{(2\pi)^3}q_x^2A(\vec{q},\omega)\Lambda(\vec{q},\omega) \qquad (19)$$

The factor of 2 is from the spin of degeneracy of electrons. Following the treatment of Hansch from Eq.(6.9) to Eq(6,13) and (6.15) in Ref. [15] , we have

$$\sigma_{xx} = \frac{e^2 n_0}{m^*} \int_0^\infty d\omega \left(-\frac{\partial n_F}{\partial \omega}\right) \frac{\Lambda(\omega)}{\Gamma(\omega)} \frac{1}{1+h^2(\omega)}$$

$$\sigma_{xy} = \frac{e^2 n_0}{m^*} \int_0^\infty d\omega \left(-\frac{\partial n_F}{\partial \omega}\right) \frac{\Lambda(\omega)}{\Gamma(\omega)} \frac{h(\omega)}{1+h^2(\omega)} \quad (20)$$

where $\Gamma(\omega)$ is the imaginary part of self-energy given by [16]

$$\Gamma(\omega) = -Im\Sigma^r(\omega) = \pi \int_{\omega_c}^{\omega_{FD}} du \alpha^2 F(u) [2n_B(\omega) + n_F(u+\omega) + n_F(u-\omega)] \quad (21)$$

The in-plane resistivity ρ_{xx} and the Hall angle $cot\theta_H$ are then

$$\rho_{xx} = \frac{\sigma_{xx}}{\sigma_{xx}^2 + \sigma_{xy}^2}$$

$$cot\theta_H = \frac{\sigma_{xx}}{\sigma_{xy}} \quad (22)$$

respectively.

Considering a high-temperature expansion of Eqs(10), (20)-(21) and assuming that u/kT is small compared to one, we may expect a linear temperature dependence of the reciprocal of longitudinal conductivity $1/\sigma_{xx}$, and the quadratic temperature dependence of the reciprocal of the Hall conductivity $1/\sigma_{xy}$, which return to the results derived from the conventional Boltzmann transport equations. However, we should remember that this is under the high-temperature conditions, $u/kT \ll 1$ or kT $\gg \omega_{FD} > 1000K$. Therefore, in the intermediate temperature range, it is difficult to get an analytical solution of the integral equation, we have to perform a numerical calculation.

In the integral of σ_{xx} and σ_{xy}, the infinity can be treated as a large but limited value ω_{HI} while the energy of electrons has a large deviation from the Fermi energy ε_F

(ε_F is set to zero). Since the terminal values of conductivities σ_{xx} and σ_{xy} controlled by the factor of $n_F(\omega_{HI})$ will be very small compared to the limited part of the integrals.

For performing the numerical calculation we need some other parameters: the Fermi wave vector $k_F = 0.35A^{-1}$ [17], the effective mass of an electron $m^* = 5m_e$, and correlation length $\xi = 20\mathring{A}$ [18] . For simplicity, we choose a coefficient in Eq.(11) to make the function $\alpha^2 F(u)$ normalized and the function $\alpha_{tr}^2(u)$ is then equal to 0.2 or so. The calculated results of the in-plane resistivity ρ_{xx}, the Hall conductivity σ_{xy} and the Hall angle $cot\theta_H$ are shown in Figs. 1-3 with the parameters, fracton Debye frequency $\omega_{FD} = 2000K$, fracton crossover frequency $\omega_c = 50K$, and the maximum value of energy in the numerical integrals (18) $\omega_{HI} = 10000K$, with different external magnetic fields. From the Figs. 1-3, we can see that the ρ_{xx} vs T curve is still nearly linear as we calculated previously. The cubed temperature dependence of the reciprocal of the Hall conductivity $1/\sigma_{xy}$ and the quadratic temperature of the Hall angle $cot\theta_H$ are well agreement with the measurements of high-T_c cuprates in the normal states. However, if different parameters are used in the calculation, for example, assumming the ω_{HI} is a function of temperature $\omega_{HI} \propto T$, then the results will exhibits another temperature power-law.

We have also performed a similar calculation in the light of electron-phonon interaction under Debye approximation. It is difficult to find the conditions so that ρ_{xx} varies T and σ_{xy} varies T^{-3} simultaneously. The results reflect that the existence of fractons essentially affect the physical properties of condensed matter. The fractons are superlocalized with a special dispersion relation and they are a kind of high-energy excitations, while the phonons are extended with frequencies less than ω_c. As Nakayama *et al.* indicated [19], in ordered structures anharmonicity reduces thermal transport but in random structures, anharmonicity is the cause of heat flow, which can account for the thermal conductivity anomalies of amorphous solids at low temperatures. For a two-dimensional percolation network, we have demonstrated that the superconductivity function $\alpha^2F(u)$ varying $u^{-1.298}$ for fractons (it varies u for phonons) may result in an increase in T_c in fractal superconductors [11]. The superconductivity function $\alpha^2F(u)$ and the transport spectral function $\alpha_{tr}^2F(u)$ of fractons differing substantially from those of phonons will also lead to the differences of their electrical transport properties.

In conclusion, we have calculated the resistivity and the Hall conductivity via temperature in the light of the fracton-electron interaction by solving the quantum transport equations. The results indicate that, under some conditions, the resistivity from the contribution of the fracton-electron interaction is nearly linear to temperature and the Hall angle $cot\theta_H$ varies T^2, which is qualitatively in agreement with the experimental results in the normal state of high-T_c cuprates. The electron-fracton interaction mechanism may offer a possible explanation to the electrical transport properties of high-T_c cuprates in the normal states.

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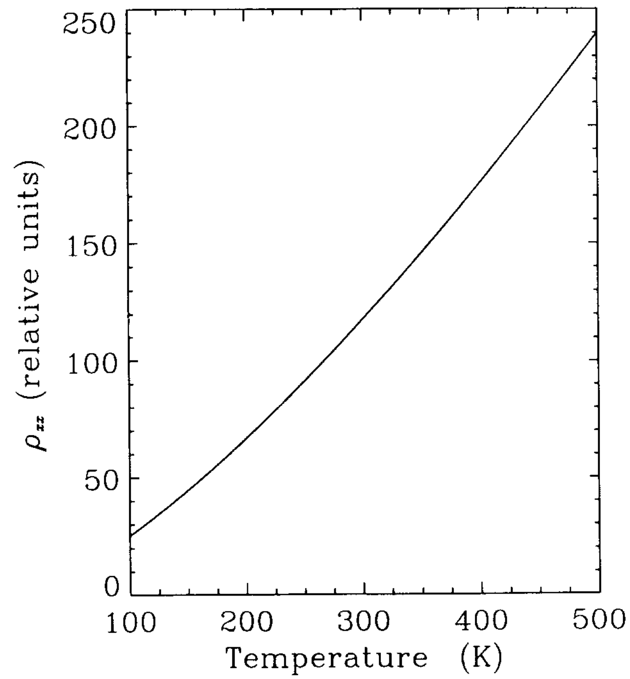

Figure 1: Temperature dependence of the in-plane resistivity ρ_{xx} with the parameters ω_c=50K, ω_{FD}=2.0$\times10^3K$ and ω_{HI}=1.0$\times10^4K$.

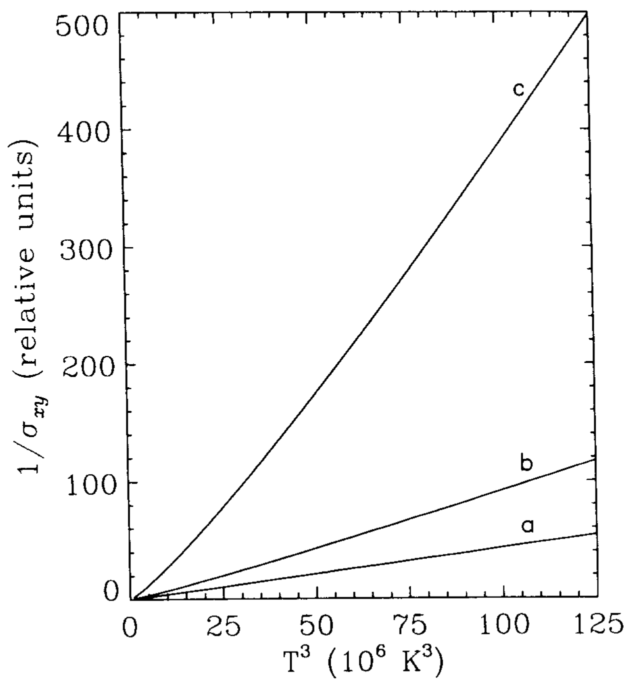

Figure 2: Reciprocal of the Hall conductivity $1/\sigma_{xy}$ as a function of temperature cubed T^3 with the parameters of ω_c=50K, ω_{FD}=2.0×10^3K and ω_{HI}=1.0×10^4K and H=8T for curve (a), H=4T for curve (b) and H=1T for curve (c).

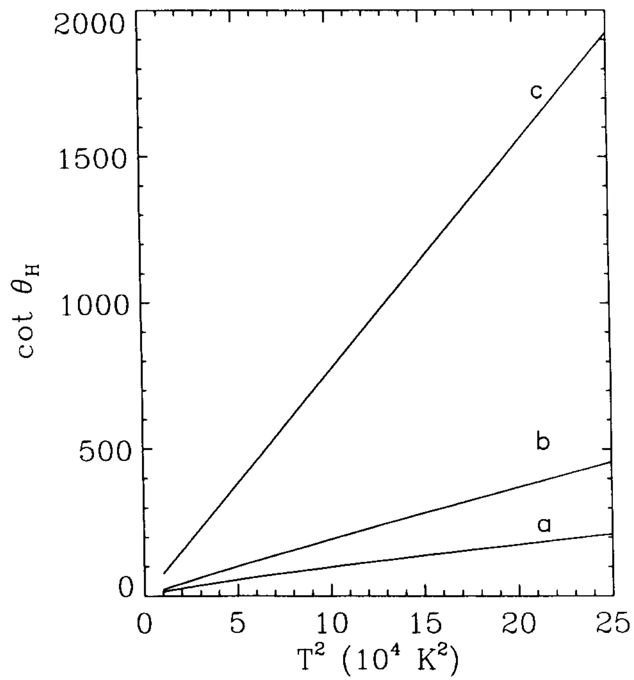

Figure 3: Hall angle $cot\theta_H$ as a function of temperature squared T^2 with the parameters of ω_c=50K, ω_{FD}=2.0×10^3K and ω_{HI}=1.0×10^4K and H=8T for curve (a), H=4T for curve (b) and H=1T for curve (c).