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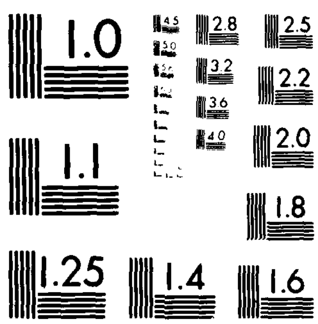

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A Scaling Theory of the Hall Effect in Disordered Electronic Systems

Boris Shapiro*†
Joseph Henry Laboratories of Physics
Princeton University
Princeton, New Jersey
08544

Elihu Abrahams‡
Serin Physics Laboratory
Rutgers University
Piscataway, New Jersey
08854

ABSTRACT

A scaling theory for the Hall effect in disordered electronic systems is developed. It is suggested that a universal scaling function for the Hall conductance exists, and the leading quantum correction to the classical value (d-2) of this function is calculated (d is the dimensionality). It is shown, by means of a scaling argument, that at the mobility edge the zero temperature Hall conductivity approaches zero with an exponent $t_H = 2t$, where t is the conductivity exponent. This relation between the exponents is supported by a microscopic calculation in $2 + \epsilon$ dimensions, which yields $t = 1$, $t_H = 2$.

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I. Introduction

The purpose of this paper is to develop a scaling hypothesis for the Hall effect in disordered electron systems. Our arguments are based on the scaling theory of electron localization which has recently been developed by Abrahams et al.[1,2] The scaling parameter in their theory is the dimensionless conductance $g(L) = G(L)/(e^2/\hbar)$ at length scale L. Here $G(L)$ is the zero temperature conductance of a hypercubic sample of size L in d dimensions. The conductance at scale bL (b is the scaling factor) is determined by the scaling relation[1]

$$g(bL) = f(b;g(L)), \tag{1.1}$$

or in continuous terms

$$d \ln g/d \ln L = \beta(g) \tag{1.2}$$

where, for each d, β is a universal function of g only.

The scaling theory predicts $d=2$ as the lower critical dimensionality in the following sense: The existence of a mobility edge is indicated by a zero of the β-function. At $d = 2$, $\beta(g) < 0$ for all finite g, and no zero occurs. For $d > 2$, $\beta(g)$ has a zero, and the conductivity

$$\sigma(E) = (e^2/\hbar) \lim_{L\to\infty} L^{2-d} g(L;E) \tag{1.3}$$

is finite for Fermi energies E higher than the mobility edge E_c. Near (above) E_c, we have

$$\sigma(E) \sim (E-E_c)^t . \tag{1.4}$$

The conductivity exponent t is related to the correlation length exponent ν by

$$t = (d-2)\nu , \tag{1.5}$$

a result first obtained by Wegner.[3]

We want to extend the scaling arguments of Ref. 1 to the Hall effect. The influence of magnetic field on the localization picture at $d = 2$ has been discussed from the microscopic point of view by several authors.[4-6] The calculations in these papers are restricted to the weak scattering limit. Here we shall give a scaling hypothesis which gives the behavior near the mobility edge. While our arguments are given for temperature $T = 0$, we expect that at finite T the length scale is set by the temperature-dependent Thouless length[7,8] $L_T \simeq (\ell_{e\ell}\ell_{in})^{\frac{1}{2}}$ where the ℓ's are the elastic and (temperature-dependent) inelastic mean-free paths.

In Section II, we formulate the scaling hypothesis. The hypothesis is supported to some extent by a microscopic calculation in Section III. In Section IV, we give an explicit calculation in $2 + \varepsilon$ dimensions where the mobility edge is accessible by perturbation theory. The conclusions are summarized in Section V.

II. The Scaling Hypothesis for the Hall Conductance

The Hall conductance G_H is defined in terms of the transverse Hall voltage U_H, the longitudinal voltage U and the conductance G by

$$G_H = GU_H/U \quad . \tag{2.1}$$

The Hall conductivity is defined as

$$\sigma_H(E) = \lim_{L\to\infty} L^{2-d}\, G_H(L;E) \; . \tag{2.2}$$

For $E < E_c$, i.e. in the insulating region, $\sigma_H = 0$ since there can be no Hall voltage without an ohmic current. When the mobility edge is approached from above, $\sigma_H(E)$ presumably approaches zero according to

$$\sigma_H(E) \sim (E-E_c)^{t_H} \tag{2.3}$$

which defines the Hall conducitivity exponent t_H.

Let us introduce a dimensionless Hall conductance $g_H(L) = G_H(L)/(e^2/\hbar)$ and try to understand how it might scale with the sample size L.

It is instructive to start with the classical transport regime. In this regime macroscopic transport theory is valid, the parameter $g(L)$ is much larger than unity and it scales with L as L^{d-2}. The Hall field $E_H = rBj$, where r is the Hall constant of the material and j is the current density. Since $U_H = E_H L$ and $jL^{d-1} = UG$, we have

$$U_H = (e^2/\hbar) rBL^{-(d-2)} gU \tag{2.4}$$

or, using Eq. (2.1),

$$g_H(L) = (e^2/\hbar) rBL^{-(d-2)} g^2(L) \quad . \tag{2.5}$$

In this classical regime $g(bL) = b^{d-2} g(L)$ and thus Eq. (2.5) implies the following scaling relation for $g_H(L)$:

$$g_H(bL) = (e^2/\hbar) rBL^{-(d-2)}\, b^{d-2}\, g^2(L) \equiv$$

$$\equiv h_{cl}^{-1}(L) b^{d-2}\, g^2(L) \quad . \tag{2.6}$$

The parameter h_{cl} has the meaning of a conductance L^{d-2}/rB, measured in units $(e^2/\hbar)$. In the classical transport regime both $h_{cl}(L)$ and the scaling parameter $g(L)$ scale as L^{d-2}, and so does the Hall conductance $g_H(L)$.

We now assume that $g_H(L)$ possesses scaling behaviour not only in the classical limit (i.e. for $g \to \infty$) but for any value of the scaling parameter $g(L)$. With Eq. (1.1) in mind, the generalization of Eq. (2.6) suggests the following scaling relation:

$$g_H(bL) = h^{-1}(L) F(b; g(L)), \tag{2.7}$$

where $h(L)$ is a (dimensionless) conductance inversely proportional to the magnetic field B.

According to the universality argument of Ref. 1 the dimensionless conductance $g(L)$ scales classically, i.e. as L^{d-2}, whenever it is large. We now assume that the same is true for $h(L)$, i.e. $h(L) \sim L^{d-2}$ for $h >> 1$. For small enough B the condition $h >> 1$ will be satisfied whatever the value of g is. In fact our assumption is that even when quantum corrections to $g(L)$ become important, $h(L)$ still scales classically if it is large enough. Fukuyama's microscopic calculation[4] in two dimensions gives some evidence in favor of this assumption as does the $2 + \varepsilon$ dimension calculation in Section III. Thus, if we require h to be large even at some typical microscopic scale L_0, we get a rough criterion for a weak magnetic field. This requirement gives $B << \hbar L_0^{d-2}/e^2 r$. (Estimating $r \approx 1/enc \approx L_0^d/ec$, we have $B << \hbar c/eL_0^2$. If L_0 is of the order of interatomic spacing this gives the usual requirement for a "classical" magnetic field in metals.)

If the conjecture about classical scaling of the parameter $h(L)$, for any value of $g(L)$, is correct, we have:

$$g_H(bL) \sim L^{-(d-2)} F(b;g(L)) \ . \tag{2.8}$$

It follows from this equation that the Hall conductivity exponent (Eq.(2.3))

$$t_H = 2(d-2)\nu = 2t \ . \tag{2.9}$$

To derive this result we introduce

$$\Delta(L) = (g(L) - g_c)/g_c \tag{2.10}$$

as the basic scaling parameter (rather than $g(L)$ itself as in Ref. 1). For L of the order of the correlation length ξ this parameter is of order unity.[9] For $L << \xi$, $\Delta(L) << 1$. In this case $\xi \sim L\Delta^{-\nu}(L)$ and thus, for $\Delta << 1$, the parameter $\Delta(L)$ scales as

$$\Delta(L) = \Delta_0(L/L_0)^{1/\nu} \qquad (2.11)$$

where $\Delta_0 \sim (E-E_c)/E_c$ is the initial value of the parameter at some scale, e.g. the microscopic scale L_0 .

We consider now a sample of size $\mathcal{L}$ and divide it into blocks of size L , i.e. $b = \mathcal{L}/L$. For a large ($\mathcal{L} >> \xi$) sample with E above E_c (i.e. in the metallic region), the Hall conductance $G_H(\mathcal{L})$ must be proportional to $\mathcal{L}^{d-2}$, which via Eq. (2.8) implies

$$g_H(\mathcal{L}) \sim L^{-(d-2)}(\mathcal{L}/L)^{d-2}\ \phi(\Delta(L)) \ . \qquad (2.12)$$

Since $\Delta \propto E - E_c$, we require, by Eq. (2.3) that for $\Delta << 1$, $\phi(\Delta) \sim \Delta^{t_H}$. Hence using Eq. (2.11) we have

$$g_H(\mathcal{L}) \sim L^{-(d-2)}(\mathcal{L}/L)^{d-2}\ \Delta_0^{t_H}\ L^{t_H/\nu} \qquad (2.13)$$

Since the block size is arbitrary (the only condition being $L << \xi$) , it must cancel out from Eq. (2.13), which immediately leads to relation (2.9) for the exponents.[10] Clearly, Eq. (1.5) for the conductivity exponent t can also be derived by a similar argument.

Taking $b = 1 + \delta$ ($\delta \to 0$) , the scaling relation (2.8) (or (2.7)) can be cast into differential form

$$d\ \ln g_H(L)/d\ \ln L = \gamma(g(L)) \qquad (2.14)$$

where $\gamma(g)$ is a universal function of the scaling parameter $g(L)$. For $g \to \infty$, i.e. in the classical transport regime, $g_H(L) \sim L^{d-2}$ and hence $\gamma = d - 2$. On the other hand, for $g \to 0$, i.e. in the strongly localized regime, $g_H(L)$, as well as $g(L)$, is exponentially small and $\gamma \to -\infty$.

Thus the qualitative behaviour of the γ-function is similar to that of the β-function.[1] However, quantitatively these two functions are different as is shown in the next section. It is this difference which accounts for the difference in the critical behaviour of $\sigma(E)$ and $\sigma_H(E)$.

Finally we would like to comment on the following point: In the above arguments, as well as in the following calculations in Sections III and IV, we ignore the dependence of the scaling parameter $g(L)$ on the magnetic field B. This is justified[6] only for $B \ll \hbar c/4eL^2$, where L is either the sample size or the Thouless length, whichever is smaller. In particular, if $T = 0$ and the sample size $L \to \infty$ (it is under these conditions when the critical exponents can be rigorously defined) any finite field would change $g(L)$ in an essential way, and thus our results for this case refer strictly speaking to an infinitesimal B. However, all the arguments of this section can be repeated without change for a finite B (which satisfies the criterion $B \ll hL_0^{d-2}/e^2r$ derived above) if we assume that the magnetic field does not represent a relevant scaling field for the Anderson transition ($d>2$), i.e. we still have a single parameter scaling theory. The fact that the scaling parameter $g(L)$ becomes dependent on B changes nothing in our arguments. Thus, if B is not relevant, we expect that the critical exponents calculated in Section IV for an infinitesimal magnetic field (since the field dependence of g is neglected) will remain unchanged at finite fields.

III. The Weak Scattering Regime

In this section we calculate the leading (i.e. proportional to $1/g$) quantum correction to the classical limit $(d-2)$ of the γ-function. Thus we consider the weak scattering regime $k\ell \gg 1$, where k is the Fermi wave number and ℓ is the electron mean-free path.

For the β-function this leading correction has been calculated in Refs. 2,11. This correction is due to the maximally crossed diagrams, which can be summed explicitly. The sample size L enters the calculation of Ref. 2 via the lower cutoff $1/L$ in some integrals over the momentum space. This leads to an L-dependent conductivity $\sigma(L)$ from which one obtains the conductance as $G(L) = L^{d-2}\,\sigma(L)$. Although the explicit calculation in Ref. 2 has been done for $d = 2$, it is trivially generalized to any dimension, with the following result:

$$g(L;E) = \frac{\hbar}{e^2}\,\sigma_0(E)L^{d-2} - \frac{A}{d-2}\,\left[\left(\frac{L}{\ell(E)}\right)^{d-2} - 1\right] \qquad (3.1)$$

Here $A = 2\pi^{-1}(2\pi)^{-d}\,S_d$, where $S_d = 2\pi^{d/2}[\Gamma(d/2)]^{-1}$ is the area of a d-dimensional sphere of unit radius. The first term in Eq. (3.1) represents the classical conductance, with

$$\sigma_0 = 2(e^2/\hbar)d^{-1}(2\pi)^{-d}\,S_d k^{d-1}\,\ell. \qquad (3.2)$$

The second term in Eq. (3.1) represents the leading quantum correction. It was assumed in the derivation of Eq. (3.1) that the sample size L is bigger than the mean-free path ℓ. The factors 2 in the expressions for A and σ_0 account for spin-degeneracy.

In the limit $d \to 2$ Eq. (3.1) takes the form

$$g = (\hbar/e^2)\sigma_0 - (1/\pi^2)\ln(L/\ell)\,, \qquad (3.3)$$

and the result of Ref. 2 is recovered.

For the perturbative calculation above to be valid the second term in Eq. (3.1) must be much smaller than the first one, which, with the help of Eq. (3.2), leads to the following criterion:

$$\alpha(L;E) \equiv \frac{d}{\pi}\,(k\ell)^{1-d}\,\frac{1}{d-2}\,[1-(\frac{\ell}{L})^{d-2}] \ll 1 \quad . \tag{3.4}$$

In three dimensions the weak scattering condition $k\ell \gg 1$ itself insures the fulfillment of criterion (3.4), for any L (bigger than ℓ) . However, near two dimensions, i.e. for $(d-2) \ll 1$, the condition (3.4) is satisfied, for any L , only if a more restrictive requirement on $k\ell$ is imposed, namely $k\ell \gg 1/(d-2)$. (As we shall see below this corresponds to energies much higher than the mobility edge.) Otherwise the condition (3.4) can be satisfied only for not-too-long samples. In particular at two dimensions the condition (3.4) reduces to $\ell n(L/\ell) \ll k\ell$, and hence for large enough L the perturbative calculation breaks down, however high is the energy. In terms of the parameter α Eq. (3.1) can be rewritten as

$$g(L;E) = (\hbar/e^2)\sigma_0(E)L^{d-2}[1-\alpha(L;E)] \quad . \tag{3.5}$$

The β-function calculated from Eq. (3.1) is

$$\beta(g) = d - 2 - A/g \tag{3.6}$$

Since A is a constant, depending only on dimensionality, the perturbation calculation supports the existence of a universal scaling function $\beta(g)$.

We now discuss quantum corrections to the Hall conducatnce $g_H(L)$. The contribution of the maximally crossed diagrams to the Hall conductivity in two dimensions has been calculated by Fukuyama[4] and, using a somewhat different technique, by Altshuler et al.[6] Again the restriction to $d = 2$ in Refs. 4,6 is not essential and in fact the value of d is introduced only at the final stage of the calculation. Employing the technique of Ref. 6 we find

$$\begin{aligned} g_H(L;E) &= \omega_c\tau\{\frac{\hbar}{e^2}\,\sigma_0 L^{d-2} - \frac{2A}{d-2}\,[(\frac{L}{\ell})^{d-2} - 1]\} = \\ &= \omega_c\tau(\hbar/e^2)\sigma_0 L^{d-2}[1-2\alpha(L;E)] \quad , \end{aligned} \tag{3.7}$$

where $\omega_c = eB/mc$ is the cyclotron frequency (m is the electron effective mass). Thus the quantum correction to the Hall conductance, relative to the leading classical term, is twice as big as the correction to the ohmic conductance (Eq. (3.5)). This result which for $d = 2$ has been derived in Refs. 4,6, holds for any dimensionality.

Differentiating Eq. (3.7) and using Eq. (3.1) for g, we obtain

$$d \ln g_H / d \ln L = d - 2 - (2A/g)(1-\alpha)/(1-2\alpha) \quad . \tag{3.8}$$

With the approximation involved in calculation of g and g_H (Eqs. (3.5), (3.7)) it would be inconsistent to keep the small α-terms in Eq. (3.8). Our approximation enables us to derive only the leading quantum correction in the γ-function:

$$\gamma(g) = d - 2 - 2A/g = -(d-2) + 2\beta(g) \tag{3.9}$$

In order to check if there are indeed no nonuniversal terms of order $\alpha A/g$ in the γ-function one needs to calculate both g and g_H to the accuracy α^2. The existence of the universal β-function implies that there is no α^2 term (nor higher order terms in α) in Eq. (3.5). For $d = 2$ the absence of the α^2 term has been proven in Ref. 11). On the contrary, if a universal γ-function exists, one must expect a term α^2 (and no higher order terms in α) in the square brackets in Eq. (3.7). This would insure that Eq. (3.9) holds also to higher order in α.

Integrating Eq. (3.9) we have

$$g_H(L) = g_{H0}(L/L_0)^{-(d-2)}[g(L)/g_0]^2 \tag{3.10}$$

where g_0 and g_{H0} are the initial values of $g(L)$ and $g_H(L)$ at some (microscopic) scale L_0. Eq. (3.10) implies that the leading quantum correction does not spoil the classical, i.e. as L^{d-2}, scaling of the parameter h(L) (Eqs. (2.7), (2.8)). This is in agreement with the assumption we made in Section 2. On this basis, we conclude that the Hall conductance exponent has the value $t_H = 2t = 2(d-2)\nu$.

4. ε - Expansion near Two Dimensions

In this section we shall be interested in the critical behaviour of $\sigma(E)$ and $\sigma_H(E)$ near the mobility edge E_c . The only physical dimension of interest in this problem is $d = 3$, since at $d = 1,2$ all the states are localized and there is no mobility edge at all.[1] Unfortunately at $d = 3$ the calculation of the preceding section are not valid near E_c , since the parameter $k\ell$ there is of order unity, while the weak scattering regime requires $k\ell >> 1$. However, at $d = 2 + \varepsilon$ with $\varepsilon << 1$, one can establish a connection between the weak scattering regime and the critical regime. This is possible because $E_c \to \infty$ when $\varepsilon \to 0$, and thus for small ε the parameter $k\ell$ remains large even at the mobility edge. In terms of the scaling parameter g this means that for small ε , g_c is large, and therefore the necessary weak scattering condition $g >> 1$ is fulfilled at the mobility edge (while for $d = 3$, $g_c \simeq 1$).

The calculation is straightforward. The β-function for small ε is given by

$$\beta(g) \equiv d\,\ell n\, g/d\,\ell n\, L = \varepsilon - 1/\pi^2 g \quad . \qquad (4.1)$$

Small terms of the order ε/g are omitted in Eq. (4.1). The zero of the β-function is $g_c = 1/\pi^2\varepsilon >> 1$. Integrating Eq. (4.1), with $g(L_0) = g_0$ as an initial condition, we obtain:

$$g(L) = g_c[1 + \Delta_0(L/L_0)^{\varepsilon}] \qquad (4.2)$$

where $\Delta_0 = (g_0-g_c)/g_c$. In fact, this equation, with a properly chosen

g_0 , is the same as Eq. (3.1). The difference is that Eq. (4.2) is valid in a much larger region than the initial Eq. (3.1). Namely, in Eq. (4.2) there is no restriction on L due to the condition $\alpha << 1$ (Eq. (3.4)). This is of course because we are relying on the universal character of the β-function. Thus the only condition for validity of Eq. (4.2) is $g >> 1$. This means that for $\Delta_0 \geq 0$ (i.e. above or at the mobility edge) Eq. (4.2) is valid for any L . On the contrary, for $\Delta_0 < 0$, i.e. below the mobility edge, Eq. (4.2) is valid for not-to-large a sample.

The exponents t and ν are immediately obtained from Eq. (4.2). For $\Delta_0 > 0$ and in the limit $L \to \infty$, g(L) is proportional to $\Delta_0 L^{\varepsilon}$. Since $\Delta_0 \sim E - E_c$, this means that the conductivity exponent $t = 1$. It follows then from the scaling relation (1.5) that the loclization length exponent $\nu = 1/\varepsilon$. The values for these exponents have been also obtained by an approach based on a Lagrangean formulation of the localization problem (see Ref. 12 and references therein).

The Hall conductivity exponent t_H is obtained in a similar way. Eq. (3.9) for the γ-function and hence Eq. (3.10) for g_H are valid if $g >> 1$. Thus, for $\varepsilon << 1$, Eq. (3.10) is valid for a sample of any size L all the way down to the mobility edge. Substituting expres.ion (4.2) for g(L) into Eq. (3.10) we obtain

$$g_H(L) = (g_{H0}/g_0^2)(L/L_0)^{-\varepsilon}\, g_c^2[1 + \Delta_0(L/L_0)^{\varepsilon}]^2 \quad . \qquad (4.3)$$

For fixed L_0 , the initial parameters g_0 and g_{H0} depend on energy , i.e. on Δ_0 . Since we are interested in the region near the mobility edge ($\Delta_0 << 1$) , we can take in Eq. (4.3) the values of these parameters at the mobility edge. It follows from Eq. (4.3) that in the limit $L \to \infty$, $g_H(L)$ is proportional to $\Delta_0^2 L^{\varepsilon}$, and thus the exponent $t_H = 2$. This result confirms the scaling relation (2.9) between the exponents.

In addition to the difference in the critical behaviour of $\sigma(E)$ and $\sigma_H(E)$ (in an infinite sample), there are some essential differences in the behaviour of $g(L;E)$ and $g_H(L;E)$ as functions of sample size L, for fixed Fermi energy E: i) At $E = E_c$ the conductance $g(L) = g_c$ and it is independent of L. In contrast, the Hall conductance $g_H(L)$ at the mobility edge does depend on L as $L^{-\varepsilon}$. ii) If the energy is fixed slightly above E_c (i.e. $\Delta_0 << 1$), the conductance $g(L)$ is a monotonically increasing function of L. However, the Hall conductance $g_H(L)$ first decreases with L and reaches a minimum at a scale equal to the correlation length $\xi = L_0\Delta^{-1/\varepsilon}$. Only for $L > \xi$ does $g_H(L)$ increase monotonically with L.

These differences between g (or σ) and g_H (or σ_H) arise from the quantitative difference between the scaling functions $\beta(g)$ and $\gamma(g)$, and in particular from the fact that these functions have zeros at different values of g. For small ε, the zeros of β and γ functions are at $1/\pi^2\varepsilon$ and $2/\pi^2\varepsilon$ respectively. Note that it is only the zero of the β-function which is associated with a critical point (the fixed point g_c of the recursion formula (1.2)). The zero of γ-function has no such meaning, since the scaling behaviour of g_H is driven by g.

Finally, at finite T and near the mobility edge the relevant length scale for the conductivity is set by the temperature-dependent Thouless length[7,8] L_T rather than by the correlation length ξ (see Ref. 9). Then the conductivity $\sigma(T)$ and the Hall conductivity $\sigma_H(T)$ are expected to be proportional to $L_T^{-\varepsilon}$ and $L_T^{-2\varepsilon}$ respectively (compared to $\sigma(E) \sim \xi^{-\varepsilon}$ and $\sigma_H(E) \sim \xi^{-2\varepsilon}$ at $T = 0$).

5. Conclusion

We have developed a scaling hypothesis for the Hall effect in disordered systems. It is suggested that a universal scaling function for the Hall conductance exists, and the leading quantum correction $(\sim 1/g)$ to the classical value $(d-2)$ of this function is calculated.

It follows from our theory that the critical exponent t_H for the Hall conductivity is twice the conductivity exponent t. Since the Hall coefficient r is proportional to σ_H/σ^2, we conclude that $r(E)$ approaches a constant value when the mobility edge is approached from above. On the other hand, the Hall mobility $\mu_H = cr\sigma$ approaches zero as $(E-E_c)^t$. The behaviour of the Hall coefficient is in agreement with the accepted view[13,14] that in a degenerate electron gas the "classical" expression $1/enc$ for the Hall coefficient (n is the electron concentration) remain approximately valid even in the regime $k\ell \sim 1$. The agreement with the picture of Refs. 13,14 of course fails when it comes to the Hall mobility near E_c, since the minimum metallic conductivity assumed in Refs. 13,14 does not occur in the scaling theory.

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20. ABSTRACT (Continue on reverse side if necessary and identify by block number)

A scaling theory for the Hall effect in disordered electronic systems is develo[ped]. It is suggested that a universal scaling function for the Hall conductance exis[ts] and the leading quantum correction to the classical value (d-2) of this functi[on] is calculated (d is the dimensionality). It is shown, by means of a scaling argument, that at the mobility edge the zero temperature Hall conductivity approaches zero with an exponent $t_H = 2t$, where t is the conductivity exponent. This relation between the exponents is supported by a microscopic calcula[tion] in $2 + \epsilon$ dimensions, which yields $t = 1$, $t_H = 2$.

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