MOTION OF A CURRENT PINCH IN A MAGNETIC FIELD IN SEMICONDUCTORS WITH AN S-SHAPED CURRENT-VOLTAGE CHARACTERISTIC

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The electrical properties of semiconductors with an S-shaped current-voltage characteristic in crossed magnetic and electric fields are considered. The uniform current distribution in electric fields that correspond to a negative differential resistance is unstable. It is shown that under these conditions and for a certain sample geometry a solitary current density wave (current pinch) is formed and moves with a constant velocity in a direction perpendicular to the electric and magnetic fields. The wave is stable for a given total current passing through the sample. This phenomenon leads to oscillations of the electric field in the external circuit. The variation of the current-voltage characteristic of the sample, due to movement of the current pinch, is considered. Estimates of the pinch velocity are made for various mechanisms of formation of the S-shaped current-voltage characteristic.

In systems with negative differential resistance, the homogeneous distribution of the current or the field, under definite conditions, becomes unstable[1,2]. In the case of a S-shaped current-voltage characteristic (with voltage ambiguity), the distribution of the current density in the sample becomes inhomogeneous and a current pinch is produced. The pinching of the current was experimentally observed in[3-6]. The pinch in sufficiently long samples is in a state of indifferent equilibrium relative to translation[7]. Therefore even weak perturbations should shift the pinch along the sample. It is clear that pinch motion can be caused by a magnetic field perpendicular to the current.

Unlike the motion of a plasmoid in a plasma injector, in the systems considered here the motion of the current pinch is not connected with the transport of mass and does not reduce to the trivial action of the Lorentz force, but is determined by thermomagnetic phenomena. Thus, in a system with the superheat mechanism[7], the carrier density is in general constant in the sample and the motion of the current pinch is a solitary wave of the effective temperature of the electrons. Since the electron mobility depends on the effective temperature, the current density wave is connected with such a temperature wave. In the case when the S-shaped characteristic is connected with breakdown or with a semiconductor-metal phase transition, the motion of the current pinch constitutes a solitary ionization wave.

PINCH VELOCITY

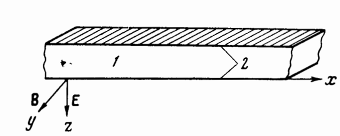
Let us consider a substance in which the negative differential resistance is the result of a sharp dependence of the conductivity on the temperature. Such a dependence $\sigma(T)$ is observed, for example, in materials with a semiconductor-metal phase transition[8], in compensated semiconductors[6], in InSb at low temperatures[9], etc. The temperature T can in this case be either the lattice temperature or the effective carrier temperature. A linear analysis of the superheat instability of a uniform distribution of the current in a transverse magnetic field B was carried out in[10]. Whereas at B = 0 the largest growth increment is possessed by waves with a wave vector **k** perpendicular to the electric field **E**, in the presence of a magnetic field the largest increment is possessed by long-wave "oblique" waves with a vector **k** making an angle

$$\varphi=\frac{\pi}{2}-\frac{1}{2}\operatorname{arctg}\frac{\sigma_{xz}'}{\sigma_{zz}'}$$

with **E**, where σ_{xz}' and σ_{zz}' are the derivatives of the components of the electric conductivity tensor with respect to temperature, **E** is directed along the z axis, and **B** along the y axis. Waves with $k \gtrsim \pi/L$ (L—characteristic thermal length) attenuate. The motion of a current pinch in a magnetic field, naturally, can be observed only in the case when the dimension of the sample in one of the directions (x, Fig. 1) greatly exceeds the pinch dimension. A linear analysis of the superheat instability for such a sample shows that if the sample dimension in the direction z is $l_z \lesssim L \tan\varphi$, then the largest increment is possessed by waves with a wave vector directed along the x axis. When this condition is satisfied, it can be assumed that the temperature and the current density do not depend on z[1)].

To determine the peculiarities of the behavior of the current pinch in a magnetic field, it suffices to consider the case when the current density and the temperature in the sample vary in the x direction. Such "one-dimensional" pinches (layers) of current

FIG. 1. Geometry of sample and its arrangement in electric and magnetic fields: 1 – semiconductor with S-shaped characteristic, 2 – metallic electrodes.



[1)]We note that when a magnetic field is applied to a sample having the large dimension along the z axis, the stationary distribution of the current may in general not have the form of a pinch, and may comprise moving "oblique" magnetic striations [11].

are stable when the sample dimension in the y direction is smaller than the pinch direction in the x direction[7].

The one-dimensional distribution of the current in the sample is determined by the following system of equations (averaged over z):

$$c\rho \frac{\partial T}{\partial t} = -\frac{\partial}{\partial x} q_x - c\rho \frac{T}{\tau} + j_z E, \tag{1}$$

$$q_x = -\varkappa \frac{\partial T}{\partial x} + \delta E, \tag{2}$$

$$j_z = \sigma(T, B) E - \beta \frac{\partial T}{\partial x}, \tag{3}$$

$$I = E l_y \int \sigma(T, B)\, dx, \tag{4}$$

where q_x is the heat flux along the sample; j_z and E are the current density and the electric field in the z direction; $B = (0, -B, 0)$; δ and β are the crossing thermomagnetic coefficients; c, κ, and ρ are the specific heat, the corresponding component of the thermal conductivity tensor, and the density; I is the total current flowing through the sample. The ambient temperature is taken to be zero. The second term in the right side of (1) describes the loss of heat from the system. The temperature τ can depend on the temperature (for example, for hot electrons). In (1)–(3), the electric-field component along the sample E_x is assumed to be equal to zero, since the sample, which is thin in the z direction, is contained between two metallic (equipotential) electrodes (Fig. 1)[2]. We note that the motion of the current pinch in the magnetic field is best observed in a sample having the form of a narrow ring, where short-circuited Hall current also flows. In the expressions for the fluxes, we disregard the terms proportional to the concentration gradients, which is valid for monopolar semiconductors with a Debye radius much smaller than L^2). In addition, it is assumed that the characteristic time of the problem L/v (v—pinch velocity) is much smaller than the Maxwellian time.

Equation (1) can be written in dimensionless variables in the form

$$\frac{\partial T}{\partial t} = \frac{\partial}{\partial x}\left(\varkappa(T) \frac{\partial T}{\partial x}\right) - \alpha(T) \frac{\partial T}{\partial x} - \varphi(T, E, B), \tag{5}$$

where

$$\alpha(T) = \frac{\tau E_1}{c\rho T_0 L}\left(\frac{\partial \delta}{\partial T} + \beta T_0\right) E, \quad \varphi(T, E, B) = T - \sigma(T, B) E^2.$$

Here the length is measured in thermal lengths $L = (\tau\kappa_0/c\rho)^{1/2}$, the time in units of τ, the electric conductivity in units of $c\rho T_0/\tau E_1^2$; as the unit of the electric field it is convenient to assume the minimum value of the field E_1, at which there still exists an ambiguity of the current-voltage characteristic (Fig. 2). The temperature and the coefficient of thermal conductivity are measured in the units T_0 and κ_0 which are characteristic of the problem. We note that the coefficient α is proportional to the magnetic field B.

An analysis of Eq. (5) in the absence of a magnetic field was carried out in[7]. Homogeneous distribution of the temperature (of the current density) corresponds to solutions of the equation $\varphi(T, E) = 0$. If the substance has an S-shaped current-voltage characteristic, this equation has three solutions in a definite region of the fields. At currents corresponding to the region of the negative differential resistance, the uniform distribution becomes unstable. This leads to the formation of the current pinch, which is stable if the total current through the sample is specified.

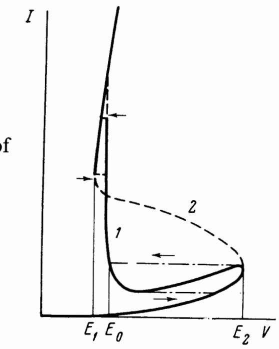


FIG. 2. Current-voltage characteristic of semiconductors; 1 – in the case of pinching of the current; 2 – in the case of uniform distribution of the current.

When the magnetic field is turned on, the current pinch (the region of increased temperature) begins to move along the sample in the direction of the action of the Lorentz force. The mechanism of this motion can be explained in the following manner. The deflection of the electrons in the magnetic field gives rise to an additional heat flux in the x direction; this flux equals $\delta(T)E$ and corresponds to the Ettingshausen flux in the case of a short-circuited current in a transverse direction. The coefficient δ for the considered materials increases with increasing temperature. This leads to a flux difference on the boundaries of the pinch. Thus, additional heating of the material occurs constantly on the leading front, and on the trailing front there is cooling, and this indeed causes the motion of the pinch. Acting in the opposite direction is the thermomagnetic current $\beta\partial T/\partial x$ (this current is the cause of the Nernst thermomagnetic effect). Indeed, on the leading front of the pinch, it leads to a decrease of the current across the sample, and consequently, to a decrease of the Joule power, and on the trailing edge it increases these quantities. This effect is always smaller than the first effect in the systems under consideration.

The solution of Eq. (5), describing the motion of the pinch, can be naturally sought in the form of a stationary solitary wave $T(x - vt)$ moving with constant velocity. Eq. (5) in a coordinate system ($\eta = x - vt$) moving together with the wave has the following form:

$$\frac{\partial}{\partial \eta}\left(\varkappa(T) \frac{\partial T}{\partial \eta}\right) - (\alpha(T) - v) \frac{\partial T}{\partial \eta} - \varphi(T, B, E') = 0, \tag{6}$$

where

$$E' = E - \frac{v}{c} B.$$

For a single hot pinch, the boundary conditions are $(\partial T/\partial \eta)|_{\eta = \pm\infty} = 0$, and the temperature T at $\eta = \pm\infty$ tends to the smaller solution of the equation $\varphi(T, B, E') = 0$.

In the absence of a magnetic field ($\alpha = 0$), the solution of Eq. (6) at specified boundary conditions exists only when $v = 0$. This can be verified by multiplying

[2]This is valid when $l_z \lesssim L$. On the other hand, if $L < l_z < d$ (d – width of pinch), then E_x differs from zero only in the wall of the pinch. The influence of this field, just as in the case when B = 0 [7], leads to renormalization of the coefficient κ, which in this case is the true coefficient of thermal conductivity.

the equation by $\kappa(T)\partial T/\partial\eta$ and integrating it from $-\infty$ to $+\infty$.

To find the velocity in a weak magnetic field B we can use the small parameter $\epsilon = \mu B/c = \Omega\tau_p$, where μ is the carrier mobility, Ω is the Larmor frequency, τ_p is the electron momentum relaxation time. From the structure of Eqs. (4)–(6) we see that the velocity v is an odd function of the field D, and the variation of the electric field an even function[3]. This makes it possible to seek the solution of (6) in the form

$$T(\eta) = T_0(\eta) + \varepsilon T_1(\eta) + \varepsilon^2 T_2(\eta) + \ldots,\quad v = \varepsilon v_0 + \varepsilon^3 v_1 + \ldots,\quad E = E^{(0)} + \varepsilon^2 E^{(1)} + \ldots. \tag{7}$$

Here $T_0(\eta) = T_0(x)$ and $E^{(0)}$ are the stationary temperature distribution and the electric field at $T = 0$. The function $T_0(\eta)$ satisfies the equation

$$\frac{\partial}{\partial\eta}\left(\varkappa(T_0)\frac{\partial T_0}{\partial\eta}\right) - \varphi(T_0, E^{(0)}) = 0, \tag{8}$$

which can be integrated in quadratures. Thus, the distribution of the temperature in the absence of a magnetic field can be regarded as known. The equation for the correction linear in ϵ is

$$\hat{L}T_1 = \frac{\alpha(T_0) - v}{\varepsilon}\frac{\partial T_0}{\partial\eta}, \tag{9}$$

where

$$\hat{L} = \frac{\partial^2}{\partial\eta^2}(\varkappa(T_0)\ldots) + \frac{\partial\varphi}{\partial T}(T_0). \tag{10}$$

The inhomogeneous equation (9) has a solution only in the case (at a value of the velocity v) when its right side is orthogonal to the solution of the adjoint homogeneous equation[12]:

$$\hat{L}^{+}\Theta \equiv \left[\varkappa(T_0)\frac{\partial^2}{\partial\eta^2} + \frac{\partial\varphi}{\partial T}(T_0, E^{(0)})\right]\Theta = 0. \tag{11}$$

We have thus obtained the condition for the determination of the velocity:

$$\int_{-\infty}^{\infty} d\eta\,\Theta(\eta)[\alpha(T_0) - v]\frac{\partial T_0}{\partial\eta} = 0. \tag{12}$$

Differentiating Eq. (8) with respect to η, we can readily verify that the function $\Theta = \kappa(T_0)\partial T_0/\partial\eta$ is a solution of (11). Since (11) is linear and the boundary conditions are homogeneous ($\partial\Theta/\partial\eta\,|_{\eta=\pm\infty} = 0$), this solution is unique. Substituting the obtained value of $\Theta(\eta)$ in (12), we obtain the value of the velocity in the approximation linear in ϵ:

$$v = \left(\int_{-\infty}^{\infty} d\eta\,\varkappa(T_0)\left(\frac{\partial T_0}{\partial\eta}\right)^2\right)^{-1}\int_{-\infty}^{\infty} d\eta\,\alpha(T_0)\varkappa(T_0)\left(\frac{\partial T_0}{\partial\eta}\right)^2. \tag{13}$$

This procedure can be continued to find the corrections to the velocity of higher order in ϵ.

We note that when the pinch moves its shape changes in the manner shown in Fig. 3a. This follows directly from the fact that the function $T_1(\eta)$, which describes the change of the shape of the pinch, is odd (this is seen from Eq. (9)).

Let us calculate the quantity $E^{(1)}$, which characterizes the change of the total resistance of the sample when the magnetic field is turned on. Since the total

[3]At a specified total current in the sample, the electric field changes when the magnetic field is turned on (owing to the motion of the pinch).

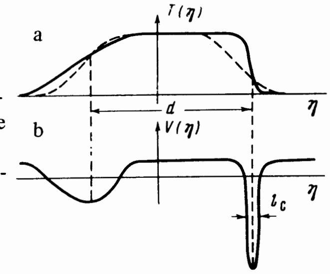


FIG. 3. Form (temperature distribution) of the moving pinch – a (the dashed line shows the distribution of the temperature in the absence of a magnetic field), and form of the "potential well" corresponding to the pinch – b.

current in the external circuit I is specified, it follows from (4) and (6) that

$$E^{(1)} = E^{(0)}\left\{1 - \left(\int\sigma(T_0)\,d\eta\right)^{-1}\int\left[\frac{1}{2}\frac{\partial^2\sigma}{\partial T^2}(T_0)T_1^2(\eta) + \frac{\partial\sigma}{\partial T}(T_0)T_2(\eta)\right]d\eta\right\}, \tag{14}$$

where $T_2(\eta)$ is determined by the equation

$$\hat{L}T_2 = \left(\frac{\alpha}{\varepsilon} - v_0\right)\frac{\partial T_1}{\partial\eta} - \frac{1}{2}\frac{\partial^2\varphi}{\partial T^2}(T_0)T_1^2 + \sigma(T_0)\left((E^{(0)})^2 - 2E^{(0)}E^{(1)}\right).$$

The quantity $E^{(1)}$ can be readily estimated for a broad pinch. Indeed, the functions $\partial\sigma/\partial T$ and $\partial^2\sigma/\partial T^2$ in (14) differ from zero only on the boundaries of the pinch (in regions of the order of the width of the wall of the pinch l_c), and therefore the ratio of the integrals in the right side of (14) amounts to a quantity on the order of l_c/d, where d is the width of the pinch (Fig. 3a). Thus, for a broad pinch $E \approx E_0(1 + \Omega^2\tau_p^2)$, i.e., the change of the total resistance of the sample in the z direction increases in exactly the same manner as the linear magnetoresistance, in proportion to $(\Omega\tau_p)^2$.

To determine the pinch velocity in a strong magnetic field ($\Omega\tau_p \gg 1$) we can use the small parameter $\epsilon' = 1/\Omega\tau_p$. The calculation procedure does not differ in general from that described above.

The method considered by us for calculating the velocity is based on the use of a small parameter connected with the value of the magnetic field. Actually we have used the smallness of the "magnetomotive force" $\alpha(T)\partial T/\partial\eta$ in (6) compared with the remaining terms of this equation. The smallness of this force is connected with the fact that it is determined by the crossing terms in the flux of (2) and (3), which in our case cannot exceed the main fluxes.

STABILITY OF MOVING PINCH

We have shown above that Eq. (5), which describes the motion of a current in a magnetic field, has a solution in the form of a solitary stationary wave moving at constant velocity along the sample. Let us ascertain the stability of such a wave.

The stability of the pinch in the absence of a magnetic field was investigated in[7,13]. A current pinch is stable only when the total current through the sample is specified (i.e., in the case of a large load resistance in the external circuit). This is connected with the fact that the perturbations leading to the change of the current in the external circuit are attenuated by the power dissipation in the load resistance. This conclusion remains valid also for a pinch in the magnetic field, since the equation of the external circuit does not change on going over to a moving system of coordinates. Thus, the stability of the moving pinch must be regarded relative to the perturbations that do not change the current in the external circuit. We must therefore take into account the perturbations that are inhomogeneous also along the y axis. Linearizing Eq. (5) relative to small deviations from the stationary distribution

$$\tilde{T}(x-vt,t)-T(x-vt)=\delta T(\eta)\cos(k_y y)\exp(-\gamma t), \qquad (15)$$

we obtain

$$\frac{\partial^2}{\partial\eta^2}(\varkappa\delta T)-\frac{\partial}{\partial\eta}(\alpha-v)\delta T-\frac{\partial\varphi}{\partial T}\delta T=(-\gamma+k_y^2\varkappa)\delta T. \qquad (16)$$

Here κ, α, and φ are functions of η, since we have substituted in them the stationary distribution $T(\eta)$, $k_y = \pi n/l_y$, and l_y is the dimension of the sample in the y direction. The stationary wave is unstable if along the eigenvalues γ of Eq. (16) some are negative.

The substitution

$$\Theta(\eta)=\varkappa(\eta)\delta T(\eta)\exp\left\{\frac{1}{2}\int_{-\infty}^{\eta}\frac{\alpha-v}{\varkappa}d\eta\right\} \qquad (17)$$

transforms Eq. (16) into the self-adjoint equation

$$\hat{H}\Theta\equiv\left[-\frac{\partial^2}{\partial\eta^2}+V(\eta)\right]\Theta=\left(\frac{\gamma}{\varkappa}-k_y^2\right)\Theta, \qquad (18)$$

where

$$V(\eta)=\frac{1}{\varkappa}\frac{\partial\varphi}{\partial T}+\frac{1}{2}\frac{\partial}{\partial\eta}\left(\frac{\alpha-v}{\varkappa}\right)+\frac{1}{4}\left(\frac{\alpha-v}{\varkappa}\right)^2. \qquad (19)$$

We note that the substitution (17) leaves the boundary conditions of Eq. (16) ($\partial(\delta T)/\partial\eta\,|_{\eta=\pm\infty}=0$) homogeneous. Thus, the problem of finding the spectrum of the Eq. (16) has been reduced to a Sturm-Liouville problem.

Let us estimate the lower bound of the negative eigenvalues γ of (18). To this end, we use the fact that when the function $\kappa(\eta)$ in the right side of (18) is replaced by its minimal value $\kappa_0 = 1$, the eigenvalues can only decrease[12], i.e., $\gamma \geq k_y^2 + \lambda$, where λ are the eigenvalues of the operator $\hat{H}$.

It is easy to verify that

$$\Theta_0(\eta)=\varkappa(T)\frac{\partial T}{\partial\eta}\exp\left\{\frac{1}{2}\int_{-\infty}^{\eta}\frac{\alpha-v}{\varkappa}d\eta\right\} \qquad (20)$$

is an eigenfunction of the operator $\hat{H}$ with $\lambda = 0$. This is connected with the fact that this function describes the change of the temperature distribution following a small shift of the pinch, relative to which the pinch is in a state of indifferent equilibrium. The function $\Theta_0(\eta)$ has one node (i.e., $\lambda = 0$ corresponds to the first excited state), therefore, according to the oscillation theorem[12], the operator $\hat{H}$ has only one negative eigenvalue. A similar method of investigation of stability was proposed by Zel'dovich and Barenblatt[14] and was used to investigate the stability of the mains in the Gunn effect[15] and of current pinches[7].

Following[7], we used for the estimate of the principal (negative) eigenvalue the fact that (18) has the form of a one-dimensional Schrödinger equation for a particle with a potential energy $V(\eta)$. This energy has in our case the form of two asymmetrical potential wells, with a width on the order of the thickness l_w of the pinch wall, separated by a distance equal to the width of the pinch d (Fig. 3b).

The difference between our case and that considered in[7] lies in the fact that the potential wells are not identical, since the temperature distributions on the front and rear ends of the moving pinch are different (Fig. 3a). The function $\Theta(\eta)$ can be approximately regarded as a superposition of two waves $\Theta_1(\eta)$ and $\Theta_2(\eta)$ describing the motion of a particle in each of the potential wells (Fig. 3b). It is easy to verify by direct differentiation of Eq. (6) that these functions correspond to the eigenvalues $\lambda = 0$, and since these functions have no nodes, the value $\lambda = 0$ is the principal one. When account is taken of transitions between two wells, this level splits into two, and the magnitude of the splitting is determined by the overlap of the functions Θ_1 and Θ_2. Thus, the ground state, to which a symmetrical combination of the functions Θ_1 and Θ_2 corresponds, has an eigenvalue $\lambda_0 \sim -\exp(-d/l_c)$, just as in the case of a pinch at rest[7]. Since $\gamma \geq (\pi/l_y)^2 + \lambda_0$, a sufficiently broad pinch moving in a magnetic field is stable.

ESTIMATES OF THE VELOCITY

Let us examine certain mechanisms for the formation of a S-shaped current-voltage in semiconductors, and let us estimate the velocity of motion of the corresponding current pinch.

At low temperatures, the mechanisms of the relaxation of energy and momentum can lead to the formation of S-shaped current-voltage characteristics[2]. Such a characteristic was observed in n-InSb at $B = 0$[9]. In this case the electron concentration and the sample does not depend on the coordinates. The thickness of the current-pinch wall in such a system is $L \sim v_T\sqrt{\tau_\epsilon\tau_p} \sim 10^{-3}$ cm, where v_T is the electron thermal velocity and τ_ϵ and τ_p are the energy and momentum relaxation times. The crossing thermomagnetic coefficients for a nondegenerate semiconductor can be written in the form[16]

$$\beta=-\frac{en\Omega}{m^*}(1-2s)\langle\tau_p^2(\varepsilon)\rangle,\quad \delta=\frac{en\Omega}{m^*}\left(\frac{5}{2}-2s\right)T\langle\tau_p^2(\varepsilon)\rangle,$$

where $\tau_p(\epsilon) \sim \epsilon^{-s}$ (here ϵ is the electron energy). Substituting these coefficients in formulas (5) and (13), we obtain the following value of the velocity:

$$v\approx A\Omega\tau_p\mu E\ [\text{cm/sec}], \qquad (21)$$

where $A = 1 + (5/2 - 2s)(1 - 4s/3)$, and the parameter s for semiconductors with S-shaped characteristic is always negative. For example, in scattering of electrons by ionized impurities $s = -3/2$, i.e., $A = 17.5$. Thus, the velocity of the pinch in this material can reach 10^7 cm/sec already at $B \sim 100$ G ($\Omega\tau_p \sim 0.1$) ($\mu \sim 5\times10^5$ cm/sec, $E \sim 1$ V/cm).

Generally speaking, in strong magnetic fields ($\Omega\tau_p > 1$), the S-shaped characteristic in a regime with a short-circuited hole current (at our sample geometry) can vanish and can even change into an N-shaped one. This leads to a vanishing of the current pinch.

When $\Omega\tau_p \gg 1$, the crossing thermomagnetic coefficients δ and β do not depend on the momentum scattering mechanism, meaning also on the temperature. If the electron energy and momentum scattering mechanisms are such that when $\Omega\tau_p \gg 1$ the current-voltage characteristic of the sample has an S shape, then the velocity of the pinch coincides with the drift velocity of the electrons in the electric field, $v = cE/B$.

We note that in a strong magnetic field the S-characteristic can be connected with breakdown, i.e., with ionization of the neutral donors, whose levels are split off the continuous spectrum by a strong magnetic field[17]. The foregoing estimates remain valid also for this case, if the Debye length of the carriers is much smaller than the thermal length L, as is usually the case under real conditions. The current oscillations observed in n-InSb in the case of breakdown in a magnetic field[17] may be connected with the motion (occurrence and disappearance) of a current pinch in the sample.

In some transition-metal oxides (VO, VO_2, Ti_2O_3, etc.), at definite temperatures, a semiconductor-metal phase transition is observed, accompanied by a jump of the conductivity (sometimes by several orders of magnitude)[8]. Naturally, such a sharp variation of the conductivity leads to the formation of an S-shaped current-voltage characteristic[18]. It is interesting to note that in this case the problem of the pinching of the current and of the motion of the pinch in a magnetic field can be solved exactly, making it possible to ascertain certain peculiarities in the behavior of narrow pinches. The results of the calculations lead to the following value of the velocity for a broad pinch:

$$v = \frac{2\varkappa}{(1+\varkappa)}\frac{\delta_M}{c\rho T_0}E \simeq \frac{2\varkappa}{1+\varkappa}\frac{n\cdot\varepsilon_F}{c\rho T_0}\Omega\tau_p\mu E. \tag{22}$$

All the quantities (δ_M, μ, ϵ_F) pertain here to the metallic phase, n is the electron density in the metallic phase, and T_0 is the temperature of the phase transition; κ is the ratio of the coefficients of thermal conductivity of the metallic and dielectric phases, and ϵ_F is the Fermi energy.

Narrow current pinches produced at $E_0 < E < E_2$ (Fig. 2) are more sensitive to the magnetic field. With increasing pinch dimensions, the velocity of their motion in a magnetic field and the change of the sample on the voltage increase.

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