"Oblique" field domains as illustrated by the Sasaki effect in semiconductors

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In an isotropic semiconductor with negative differential resistance, the walls of the produced strong- and weak-field domains are inclined at an oblique angle to the current direction (and not at a right angle as in the isotropic case). The theory of such oblique field domains, with the Sasaki effect in a two-valley semiconductor as the example, is developed for samples with limited transverse dimensions. The distinguishing feature of this example is that the negative resistance and the anisotropy itself are induced by heating of the carriers in the longitudinal and transverse fields and are interrelated. The inclination angles of the oblique domain walls depend on the flowing current and in one of the limiting cases they can be almost parallel to the current direction. A correct value of a strong-field domain velocity can be obtained only if the inclination is taken into account.

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1. In heating electric fields, the electric conductivity of semiconductors with cubic symmetry is no longer isotropic (the Sasaki effect). When current is made to flow through such a semiconductor with limited transverse dimensions, a transverse electric field is produced in it. If the semiconductor is characterized by a negative differential conductivity of N-type (N-NDC) and a homogeneous field distribution in it is unstable, strong- and weak-field domains are produced, in which not only the longitudinal but also the transverse electric fields are different. The existence of unequal transverse electric fields makes the standard homogeneous theory of domain production,[1,2] in which the transverse field is completely disregarded, unsuitable for this case.

We consider here field domains, using as the example the N-NDC produced in the Sasaki effect in a multivalley semiconductor. We choose for simplicity a two-valley model (Fig. 1); the axes of the large masses of the valleys (ellipsoids of revolution) are oriented at right angles to one another and lie in the xy plane, with the x axis directed along the current density far from the domain walls and the sample limited in the y direction (plate of thickness $2d$). The valley symmetry axis makes an angle ψ with the x axis.

The Sasaki effect is accompanied (and is governed to a considerable degree) by resettlement of electrons from one valley to another (the valleys are equivalent under equilibrium conditions). The difference between the heating of the electrons in the domains and their resettlement lead to the appearance of an off-diagonal conductivity component $\sigma_{xy} = \sigma_{yx}$, which increases abruptly with increasing electric field E_x if the latter is weak enough. The N-NDC considered here is due precisely to this rapid growth of σ_{xy}, since the total conductivity of the sample

$$\sigma=\sigma_{xx}-|\sigma_{xy}|^2/\sigma_{yy} \tag{1}$$

decreases rapidly in this case,and this rapid decrease of $\sigma(E_x)$ can cause a decrease of the current density $i = \sigma E_x$. This N-NDC mechanism was apparently first considered in Ref. 3, with the multiply valued Sasaki effect (MSE) as the example. It takes place at $\psi = 0$, and also in a certain range of $|\psi| < \psi_0$ near this current direction. In the MSE, the state with equilibrium populations of valleys 1 and 2 unstable at $\psi = 0$, and the other two states with unequal valley populations and with nonzero transverse field $E_y (\vartheta = E_y/E_x \neq 0)$ are stable. The state $\vartheta = 0$ is not realized in a certain range of fields $|E_x|$, starting with the value E_c (Fig. 2); in this case the corresponding branch 3 of the current-voltage characteristic (CVC) of Fig. 2 is not realized, but the degenerate branch 2,

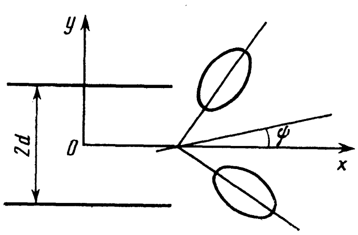

FIG. 1. Two-valley model.

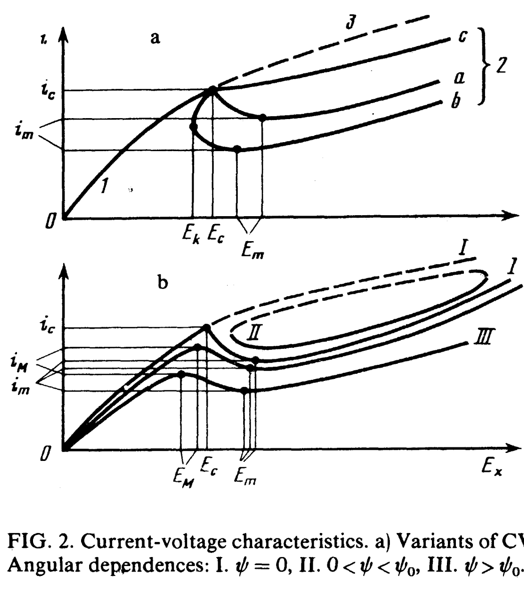

FIG. 2. Current-voltage characteristics. a) Variants of CVC at $\psi = 0$. b) Angular dependences: I. $\psi = 0$, II. $0 < \psi < \psi_0$, III. $\psi > \psi_0$.

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corresponding to anisotropic states with $\vartheta = \vartheta^{(\pm)}(E_x)$, is realized. Depending on the parameters of the electrons in the valleys, three variants of branch 2 are possible, shown by curves a, b, and c in Fig. 2a. It appears that for the MSE the usual branch is a with the N-NDC section. Branch b, which is realized in the case of very large resettlement of the electrons among the valleys, contains in the field section (E_k, E_c) two values of the current and corresponds not to two but to four anisotropic states at a given field E_x, namely $\vartheta_1^{(\pm)}$ and $\vartheta_2^{(\pm)}$. The resultant multiply valued character of the CVC is more complicated than the usual N-NDC, but from the point of view of the domain solutions investigated below the branches a and b are equivalent, so that no distinction is made between these variants of the behavior. The branch c without an N-NDC section occurs at low anisotropy of the electrons in the valleys (much lower than, for example, in germanium and in silicon), and is of no further interest to us.

When the direction of the current in the sample deviates from $\psi = 0$, the CVC undergo the changes illustrated in Fig. 2b; they are described in detail in Ref. 4. In the angle range $|\psi| < \psi_0$ where the MSE is preserved, the degeneracy of branches of Fig. 2a is lifted at $\psi \neq 0$. Besides the CVC branch corresponding to the "correct" sign of the Sasaki field E_y and containing an N-NDC, there appears a CVC loop with "incorrect" sign of E_y. In many cases the contributions of the states corresponding to the loop does not appear, and its existence can be neglected. At $|\psi| > \psi_0$ the CVC loop vanishes, but the N-NDC section is as a rule preserved. As shown by Mitin and the author (see Ref. 4), if the MSE and the associated N-NDC take place at $\psi = 0$ (Fig. 2a), and if the N-NDC mechanism due to the transfer of the electrons into the "heavy" valley is realized,[5] then N-NDC due to a combination of all these mechanisms appears in the entire range of angles ψ, and it is only in the direction $\psi = \pi/4$ and its equivalents that no transverse fields E_y are produced. We note also that the N-NDC due to the rapid growth of the off-diagonal conductivity σ_{xy} takes place at $\psi \neq 0$ also when there is no MSE at all, but a situation close to MSE obtains, i.e., the time of intervalley scattering decreases rapidly as the electron gas is heated, and the resettlement of the electrons in the valleys increases rapidly.

In the foregoing cases (except for the $\psi = \pi/4$ direction and its equivalents) it is necessary to treat correctly, when constructing a domain theory, both the longitudinal and transverse electric field. It is shown below that this can be done in "oblique"—with domain walls that are not perpendicular to the current direction, but are at oblique angles. The latter are determined by the continuity conditions of the current and of the electrostatic potential. In particular, for the inclination angle γ of a wall (Fig. 3a) between extended domains with different values of the longitudinal $[E_x^{(1)}(i)$ and $E_x^{(2)}(i)]$ and transverse $(E_y^{(1)}(i)$ and $E_y^{(2)}(i))$ fields we have

$$\tan\gamma(i) = -[E_x^{(1)}(i) - E_x^{(2)}(i)]/[E_y^{(1)}(i) - E_y\ \ (i)], \quad (2)$$

with the current density i uniform in the entire sample except in the domain wall and in surface layers with thickness on the order of that of the domain wall. At $\psi = 0$ and in the presence of the MSE, stratification of the strong-field domain is possible[6,7] into layers with different transverse field directions $\vartheta^{(+)}$ and $\vartheta^{(-)} = -\vartheta^{(+)}$. Since this is not a rough state, the interlayer wall parallel to the plate surface (Fig. 3b) should in practice always shift toward one of the sample surfaces and produce a surface wall having a complicated structure. Measurements of the anomalous Hall effect[7] by Asche *et al.*[8,9] show that although the anomalous effect due to the shift of the interlayer wall in a magnetic field does exist and is large, it is nonetheless noticeably weaker than the theoretically predicted one, apparently because of the "slowing down" of the interlayer wall by small inhomogeneities. From this point of view, the situation in Fig. 3b is of interest; since the fields E_x^0 are equal in the layers $2^{(-)}$ and $2^{(+)}$ (as are also the current densities in them), this situation reduces in practice to the case of a homogeneous strong-field domain (Fig. 3a).

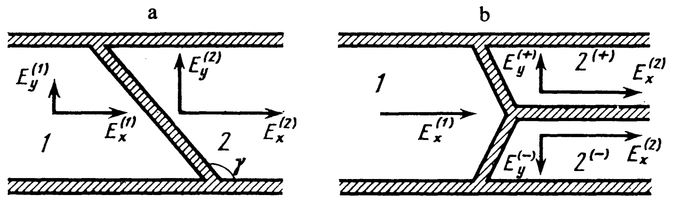

FIG. 3. Domain walls. a) Homogeneous strong-field domain ($\psi < 0$). b) Stratification in a strong-field domain.

2. We now find the quasistationary and quasineutral distributions of the electron densities $n_{1,2}$ in the valleys (these densities determine in turn the distributions of the fields $E_{x,y}$ and currents $i_{x,y}$) in an oblique domain wall (Fig. 3). It follows from the quasineutrality conditions that

$$n_{1,2} = n_0(1 \mp f), \quad (3)$$

where n_0 is the equilibrium density of the electrons in one valley. For the uniform distributions that are realized far from the domain walls, f is obtained from the condition $n_1/\tau_1 = n_2/\tau_2$ for the balance of the intervalley transitions, in the form

$$f = -(\tau_1 - \tau_2)/(\tau_1 + \tau_2), \quad (4)$$

where $\tau_{1,2}$ is the time of departure of the electrons from valleys 1 and 2 (into valleys 2 and 1, respectively).

We use the phenomenological effective-field method,[3,4] according to which it is assumed that the tensors $\hat{\mu}_{1,2}$ of the electron mobilities in the valleys and $\hat{D}_{1,2}$ of the diffusion coefficients are written in the form

$$\hat{\mu}_{1,2} = \hat{a}_{1,2}\mu(E_{1,2}), \quad \hat{D}_{1,2} = \hat{a}_{1,2}D(E_{1,2}), \quad (5)$$

where $\hat{a}_{1,2}$ are dimensionless tensors with components

$$a_{1,2xx} = 1 \pm as, \quad a_{1,2yy} = 1 \mp as, \quad (5')$$

$$a_{1,2xy} = a_{1,2yx} = \mp ac.$$

Here and elsewhere $s \equiv \sin 2\psi$ and $c \equiv \cos 2\psi$; the effective fields are given by

$$E_{1,2}^2 = \mathbf{E}\hat{a}_{1,2}\mathbf{E}. \quad (6)$$

From (3), (4), (5), and the condition $i_y = 0$ that there be no transverse current, a condition satisfied in semi-infinite domains far from the walls, we have

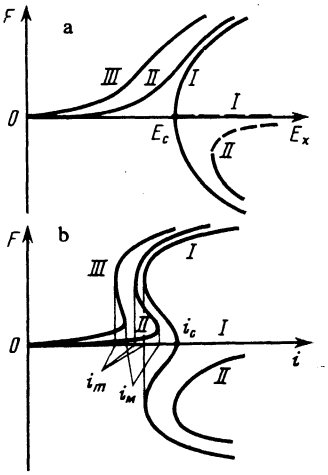

FIG. 4. Dependences of the parameter F on the longitudinal electric field (a) and on the current density (b): I, II, and III are the same as in Fig. 2.

$$\vartheta=-aFc/(1+aFs);$$
$$F=-[\Phi(E_1)-\Phi(E_2)]/[\Phi(E_1)+\Phi(E_2)], \tag{7}$$
$$\Phi(E)=\mu(E)\tau(E), \quad \tau_{1,2}=\tau(E_{1,2}).$$

The quantity F characterizes the resettlement in the valleys. Equation (7) with allowance for (6) defines the $F(E_x)$ dependence. These dependences for $\psi=0$ and for two values of $\psi>0$ (one of which is $\psi<\psi_0$ and the other $\psi>\psi_0$) are shown in Fig. 4a. Using these $F(E_x)$ dependences and the CVC $i(E_x)$ of Fig. 2b (for identical values of ψ) we easily obtain the functions $F(i)$; they are plotted in Fig. 4b (and, as can be easily verified, are qualitatively equivalent for the CVC variants a and b of Fig. 2a). For each value of the angle ψ there exists a range of currents from $i_m(\psi)$ to $i_M(\psi)$, in which one value of i corresponds to three positive values

$$F_1(i)<F_0(i)<F_2(i).$$

At $\psi=0$ the value of i_M coincides with i_c, and $F_1(i)=0$ for all i. The CVC shown in Fig. 2 follow from the formula

$$i=2eE_xn_0\frac{\Phi(E_1)+\Phi(E_2)}{\tau(E_1)+\tau(E_2)}\frac{1-a^2F^2}{1+aFs}. \tag{8}$$

The functions $f=f(i)$ have qualitatively the same form as $F(i)$, i.e., they are triple-valued in the current interval $(i_m(\psi), i_M(\psi))$ (if only $f>0$ is taken into account), since the rapid growth of the transverse field E_y in real crystals is connected precisely with the abrupt $\tau(E)$ dependence that obtains in low-temperature intervalley scattering (with emission of intervalley phonons). Neglecting everywhere below the slow and therefore qualitatively insignificant $\mu(E)$ and $D(E)$ dependences (and assuming both these quantities to be constant), we have $F=f$,

$$\vartheta=-fac/(1+fas), \tag{7'}$$

$$E_x=\lambda(1+fas)/(1-a^2f^2), \tag{8'}$$

$$E_{1,2}^2=\frac{\lambda^2}{(1-a^2f^2)^2}\times[(1+afs)^2(1\pm as)^2+(fac)^2(1\mp as)^2\pm 2fa^2c^2(1+fas)], \tag{6'}$$

where $\lambda=i/2en_0\mu$.

3. We now proceed directly to the problem of the domain wall. The latter is assumed planar, so that all the quantities depend only on the single coordinate

$$u=x+Ay. \tag{9}$$

The current-continuity equation is

$$\frac{d}{du}(j_x+Aj_y)=0,$$

whence

$$j_x+Aj_y=j\equiv i/e. \tag{10}$$

The analogous connection between the components E_x and E_y, obtained from the condition that the electric vector has no curl

$$\frac{d}{du}(E_y-AE_x)=0,$$

is of the form

$$E_y-AE_x=E_y^{(1)}-AE_x^{(1)}=-\lambda\chi, \tag{11}$$

$$\chi=[f_1ac+A(1+f_1as)]/(1-a^2f_1^2), \tag{11'}$$

We consider two limiting situations: 1) a single domain wall separating two infinite domains along the x axis (Fig. 3); 2) a thin strong-field domain ("hot" domain) or weak-field ("cold") domain against a background homogeneous along the x axis. The latter case means that the domain is thin compared with the sample thickness $2d$. Since we are considering only thick samples, much thicker than the domain walls, the thin hot or cold domains can also be substantially thicker than their own walls (and in this sense the domains need not be thin).

In the case of a single wall (Fig. 3) we have

$$E_y^{(1)}-AE_x^{(1)}=E_y^{(2)}-AE_x^{(2)},$$

whence

$$A=-\beta c/(\alpha+\beta s)\equiv A_\infty, \tag{12}$$

where $A_\infty=\cot\gamma(i)$ [cf. (2)], and

$$\chi=\chi_\infty=-c/(\alpha+\beta s). \tag{11''}$$

Here $\alpha=a(f_1+f_2)$, $\beta=1+a^2f_1f_2$; $f_{1,2}$ are the values of f in the domains. At $\psi=0$ we have $A_\infty=-1/af_2$, $\chi_\infty=-1/af_2$.

In the case of a thin domain between two semi-infinite ones in which $i_y=0$, we must impose the condition

$$\int_{-\infty}^{\infty} i_y(x)\,dx=0, \tag{13}$$

which determines A (and the angle of inclination of the entire thin domain). Since the fields $E_x^{(1)}$ and $E_y^{(1)}$ in semi-infinite domains surrounding the thin domain were obtained from the condition $i_y=0$, only the thin domain itself and its walls contribute to the integrand of (13).

When (11) is taken into account we get from (10)

$$E_x=\frac{\lambda[1+\chi(A+fac+Afas)]+(aD/\mu)\Gamma_1\partial f/\partial u}{\Gamma_2-af\Gamma_1}; \tag{14}$$
$$\Gamma_1=s(1-A^2)-2Ac, \quad \Gamma_2=1+A^2.$$

To obtain $f(u)$ we must write down and solve the continuity equation for the difference between the electron fluxes in the valleys, $\mathbf{j}'=\mathbf{j}_1-\mathbf{j}_2$ ($\mathbf{j}_{1,2}$ are the electron fluxes in the valleys); this equation takes the form

$$-\frac{\partial}{\partial u}(j_x'+Aj_y')=2\left(\frac{n_1}{\tau_1}-\frac{n_2}{\tau_2}\right)+\frac{\partial}{\partial t}(n_1-n_2). \qquad (15)$$

Substituting in (15) the values of $j'_{x,y}$ and $n_{1,2}$ obtained by using (11) and (14), and assuming that the needed solution of (15) takes the form of a stationary wave

$$f=f(u-vt)=f(x+Ay-vt), \qquad (16)$$

where v is the wave velocity along the x axis, we transform (15) into

$$Dp\frac{dp}{df}+p(v+\lambda\mu V(f))=R(f), \qquad (17)$$

where

$$p=\frac{\Gamma_2^2-a^2\Gamma_1^2}{\Gamma_2-af\Gamma_1}\frac{df}{du}, \qquad (18)$$

$$R(f)=\frac{\Gamma_2^2-a^2\Gamma_1^2}{\Gamma_2-af\Gamma_1}\frac{1}{\tau}\left(f+\frac{\tau_1-\tau_2}{\tau_1+\tau_2}\right), \qquad (19)$$

$$V(f)=-\frac{\chi\Gamma_3}{\Gamma_1}+\frac{\Gamma_2^2-a^2\Gamma_1^2}{(a\Gamma_1 f-\Gamma_2)^2}\left(1+\frac{\chi\Gamma_3}{\Gamma_1}\right),\quad \Gamma_3=c(1-A^2)+2As. \qquad (20)$$

At $\psi=\pi/4$ we have $A=0$, $\chi=0$, $\Gamma_1=\Gamma_2=1$, so that

$$\frac{\Gamma_2^2-a^2\Gamma_1^2}{\Gamma_2-af\Gamma_1}=\frac{1-a^2}{1-af},\quad V(f)=\frac{1-a^2}{(1-af)^2}. \qquad (20')$$

The singular points (17) lie on the $p=0$ axis, and at $f\geqslant 0$ their number in the current interval of interest to us is always three: $f_1\leqslant f_0\leqslant f_2$. The outer points are saddles and the central one can be a node or a focus.

4. In the case of a single static ($v=0$) domain wall, Eqs. (11) and (12) are valid. Assume that in the entire current interval $(i_m(\psi), i_M(\psi))$ they are large enough to satisfy the condition

$$\lambda^2\mu^2V^2(f_0)\gg|DR'(f_0)|, \qquad (21)$$

where $R'(f)=dR/df$. The point f_0 is then a node, and the saddles f_1 and f_2 do not have the common separatrix needed for the existence of the sought single domain wall. The situation arising is well known in domain theory[1] and can be clearly tracked at $\psi=\pi/4$, when Eq. (20') holds for $V(f)$. The necessary trajectories are obtained only in two limiting cases: at $i=i_m(\psi)$, when $f_0=f_2$, and $i=i_M(\psi)$, when $f_0=f_1$. The merging points f_0 and f_2 or f_0 and f_1 form a complicated singular point of the saddle-node type. At any ψ, two saturation currents are possible on the static CVC of a long ($l\gg 2d$) sample: $i_m(\psi)$ or $i_M(\psi)$ (Fig. 5). The choice of one of these two CVC variants is determined by the boundary conditions on the current contacts. These conditions themselves, as well as the distributions of the concentrations and of the fields at the contacts, are not considered here in view of the difficulty of the ensuing problem.

It follows from the foregoing that in the static case the problem with the oblique wall ($\psi\neq\pi/4$) differs from that of the right-angle wall ($\psi=\pi/4$) only in the very fact that the wall is slanted. When condition (21) is satisifed, we can determine $f(u)$ by using the approximate drift dependence

$$p\approx R(f)/\lambda\mu V(f).$$

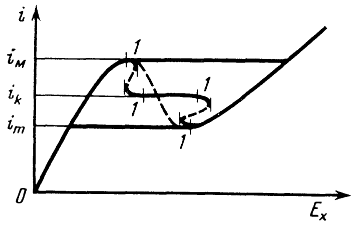

FIG. 5. CVC of sample with domain wall or with domain

5. Corresponding to a narrow hot or cold domain is the separatrix loop of the saddle f_1 or of the saddle f_2. If (21) is satisfied such a loop, which is a limit cycle, is possible under the conditions of a stationary wave with velocity close to

$$v=-\lambda\mu V(f_0). \qquad (22)$$

[By itself, rigorous satisfaction of condition (22) does not coincide with the rigorous criterion for the existence of the separatrix loop, since the point f_0 becomes not a center but only a focus of unity multiplicity.[1)] In the case of a noticeable deviation from (22), however, f_0 becomes a node; the separatrix loop exists therefore at a velocity v close to (2).]

The maximum value $f=f_M$ in a hot domain is obtained from the condition

$$\int_{f_1}^{f_M} df[V(f)-V(f_0)]\approx 0, \qquad (23)$$

which yields

$$f_M=\frac{\Gamma_2}{a\Gamma_1}-\frac{(\Gamma_2-a\Gamma_1 f_0)^2}{a\Gamma_1(\Gamma_2-a\Gamma_1 f_1)}. \qquad (24)$$

from the condition $f_M<f_2$ we obtain a constraint on the currents at which a hot domain exists:

$$(\Gamma_2-af_1\Gamma_1)(\Gamma_2-af_2\Gamma_1)<(\Gamma_2-af_0\Gamma_1)^2. \qquad (25)$$

Reversal of the inequality (25) means a transition to cold domains in a hot medium with $f=f_2$. There exists a current $i=i_k$, determined by equality of the left and right-hand sides of (25) and dividing the current interval (i_k, i_M) into two. The CVC of a long sample with a moving domain should contain a saturation section at the current $i=i_k$ (Fig. 5).

To determine all the details of the CVC we must know the parameter A obtained from the condition (13), which acquires after the calculation i_y the form

$$\oint\frac{df}{p}\frac{A(1+fas)+fac-\chi(1-a^2f^2)}{(\Gamma_2-af\Gamma_1)^2}=0; \qquad (26)$$

the integral in (26) is calculated along the separatrix loop.

As f_M approaches f_2 in the hot domain (in this case the current i decreases to the value i_k), the main contribution to the integral of (26) is made by the region with almost constant $f\approx f_2$, so that (26) reduces to the condition

$$A(1+f_2as)+f_2ac-\chi(1-a^2f_2^2)\approx 0, \qquad (27)$$

from which, taking (11′) into account we have $A=A_\infty$, $\chi=\chi_\infty$. We note that the fields $E_x^{(1,2)}(i)$ and $E_y^{(1,2)}(i)$ are different at $i=i_m$, i_M, i_k, so that the values of A_∞ and hence of the inclination angle γ are also different on all the three saturation sections shown in Fig. 5.

In narrow hot domains f_M is much smaller than f_2, so that the parameter $A(i)$ should differ noticeably from $A_\infty(i)$. We consider the onset of a hot domain when the condition

$$(f_M-f_1),(f_0-f_1)\ll f_2-f_1 \qquad (28)$$

is satisfied. In this case $f_M\approx 2f_0-f_1$, and the condition (26) reduces to

$$\int_{f_1}^{2f_0-f_1} df[aA(f+f_1)+(c+As)(1+a^2ff_1)]\approx 0, \qquad (29)$$

whence

$$A=-c[s+a(f_0+f_1)(1+a^2f_1f_0)^{-1}]^{-1}. \qquad (30)$$

The value of A given by (30) differs substantially from A_∞. As the current decreases from i_M to i_k, noticeable rotation of the hot domain takes place. The rotation effect is particularly clear at $\psi=0$, when $f_1=0$ and $A=-1/af_0$, i.e., $\tan\gamma(i)=af_0$. At $i=i_c$ we have $f_0=0$, i.e., the hot domain tends to be created along the current direction (to the extent allowed by the sample length), and only with decreasing current does it rotate towards the position given by Eq. (2).

A nascent cold domain ($f_2-f_m, f_2-f_0 \ll f_2-f_1$) can also be considered similarly and is also noticeably inclined relative to a wide domain (although this difference is not as large as in the case of a narrow hot domain at $\psi=0$). We then obtain

$$A\approx -c[s+2af_2/(1+a^2f_2^2)]^{-1}.$$

In the literally considered situation with an N-NDC stemming from the Sasaki effect it is difficult to observe the domain rotation with changing current, since the rotation situation corresponds to those sections of the CVC on Fig. 5 which are for the most part not realized in a static experiment. (In Fig. 1 the domain rotation should be observed on the sections marked with the number 1.) To increase the length of the observed sections the field section of the N-NDC must be wider than predicted and observed in germanium and silicon. This situation is realized, for example, in the case of the Gunn effect in gallium arsenide. Although the nature of the N-NDC is there quite different, in a strong-field domain a considerable fraction of the electrons stays in the L valleys, and when the direction of the current in a sample with limited transverse dimensions deviates from the symmetry axes, there should be realized there the Sasaki effect with transverse fields, oblique domains, and their rotation when the width of the domain and the maximum current in it change.

We have dealt above only with the anisotropy induced by carrier heating when current flows through the sample. It is possible to consider similarly the case with an initially anisotropic material, as well as with an isotropic material in a magnetic field. In the latter case the transverse field is the Hall-effect field, which is different in the hot and cold domains because the kinetic parameters are different there.

The author thanks O. G. Sarbeĭ for a discussion.

[1)]We recall[10] that a small change in the parameters [e.g., a change of v relative to the value (22)] transforms a focus of multiplicity 1 into a rough focus, and a limit cycle that contains this focus is produced in one of the directions of the variation. With increasing deviation of v in this direction from the value (22), the cycle broadens and at a certain value v' occupies the position of a separatrix loop. When (21) is satisfied the value of v' differs insignificantly from the value of (22), and this is why this equation can be used.

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