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Spatial Modulation of Polarization in Ferroelectric Liquid Crystals

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The formation of quasiequilibrium modulated structures with large spatial periods in the ferroelectric chiral smectics C with large magnitudes of the spontaneous polarization is discussed. The specific features of such structures are related to the nonequilibrium process of charge transfer, to the formation of orientational defects and polarization inhomogeneities stabilized by free charge carriers.

INTRODUCTION

Recently the paper[1] announced about observations of the periodic domain structure in ferroelectric liquid crystals with the period which is much larger than the period of thermodynamically equilibrium helicoidal structure in such materials and is approximately inversely proportional to the spontaneous polarization squared (at fixed other material parameters and temperature). There was observed a difference in origin conditions for the static and dynamic structures in Reference 1, where the preliminary model of distributions of molecular orientations and spontaneous polarization in a liquid-crystalline film was proposed. Authors[1] discussed the possibility of existence of a proper ferroelectricity which causes domains with wide transition zones of the diffuse Bloch wall type. The essential role of the volume electric charge screening the spontaneous polarization was noted.

These and new experimental data concerning such structures are interesting from the point of view of effect of the large spontaneous polarization, impurity charges, boundary conditions on a behaviour of ferroelectric liquid-crystalline films in external fields. The interpretation of observed phenomena demands the more detailed physical consideration of these effects. In this paper we make an attempt to explain qualitatively the mentioned phenomena to compare theoretical estimations with experimental data and to predict new effects.

It should be noted that the appearance of proper ferroelectricity in liquid crystals with the large spontaneous polarization P_c has no sufficiently serious foundations. Firstly, it is difficult to interpret the value $P_c \sim 10^{-7}$ C.cm^{-2} at relatively high temperatures $T \sim 300$K as the proper polarization magnitude in such molecular systems since the interaction of neighbouring dipoles is relatively weak. Secondly, the ordinary equilibrium helicoidal structure appears at the same temperature of

the phase transition smectic A - smectic C and so it is difficult to expect that different physical reasons cause the same phase transition point or, on the contrary, the same physical reason causes different modulated structures. Thirdly, the non-equilibrium periodic structure appears always after or during the external field action on the film of the chiral smectic phase C*. In some samples the textures mentioned above transformed in time (during tens of minutes) to ordinary equilibrium helicoidal structures with relatively small pitch of helix. Thus, in some extent, the considered structure is induced and related to definite processes following in the C* film under the field action.

In connection with these remarks it is reasonable to consider the another physical mechanism of origin of such structures related to the nonequilibrium process of charge transfer, the formation of orientational defects and polarization inhomogeneities, stabilized by free charge carriers.

THE NONEQUILIBRIUM STATE OF A LIQUID-CRYSTALLINE FERROELECTRIC AFTER SWITCHING OFF THE EXTERNAL ELECTRIC FIELD

The behaviour of ferroelectric liquid crystals after switching on and switching off the external field E is not enough investigated in the situation when such a ferroelectric has the large spontaneous polarization and considerable quantity of free electric charge carriers. In absence of the external field the polarizing role of the boundary surfaces of the sample with the relatively large thickness $d \sim 10 \div 30$ mcm is weak if the screening radius r_D is small. In this case the appearance of smectic phase at the temperature change must be accompanied by the spontaneous appearance of the orientation helix $(n_x(z), n_y(z))$ and the polarization helix $(P_x(z), P_y(z))$ with the thermodynamically equilibrium pitch h_0 (the wave number $q_0 = 2\pi/h_0$) which is slightly perturbed by the subsurface electric field E_s existing in the thin double layer with the thickness r_D.

The action of the sufficiently strong external field E causes untwisting of the polarization helix $(P_x(z), P_y(z)) = P(\cos\varphi, \sin\varphi)$ and the emergence of excess of ions with different signs at the opposite boundaries of the sample. At the following switching off the external field the volume charge starts to bleed under the action of the field E' occurred through the surface charge density and the process of recovery of the orientation helix starts too. One can consider the formation of helix as the process of appearance of walls with the thickness η, the orientations of director n and polarization P sharply change inside these walls. In presence of the field E' the wall thickness is $\eta \sim (K\theta/\mu E')^{1/2}$,[2] where K is the modulus of orientational elasticity, θ is the tilt angle in the C* phase, μ is the piezoelectric modulus. If the field E' is smaller than the certain critical value E_c, the quantity of walls is large and the magnitude of the pitch of helix is close to the thermodynamically equilibrium value. However, such a situation takes place if the director n and the polarization P weakly interact with sample boundaries.

If in the beginning of the relaxation process, when the screening is weak, the subsurface field is strong, the polarization P remains fast coupled with a surface,

for example P is parallel with the normal to surface. In this case the initial helix cannot restore spontaneously since, for topological reasons, the strict orientation at the sample boundaries inevitably results in appearance of orientation and polarization defects (disclinations) near by the sample surfaces.[3] The disclinations, precisely the disclinational loops, and walls mentioned above play, in fact, the same role in the helix formation, this helix can be pictured as the periodic system of extensive disclinational loops and walls connected with loops.[4] During the relaxation process initially the quantity of such loops is small since their origin demands the energetic expenditure, the field E' being still strong.

Thus the start of relaxation after switching off the external field must be accompanied by the origin of a helix with the relatively large pitch $h >> h_0$. In this case flowing of the volume charge can play the essential role for the stabilization of such a helix which is thermodynamically nonequilibrium and has the intermediate value of pitch h. This stabilization is conditioned by the following reasons: there is an inhomogeneity of the macroscopic polarization in the vicinity of the disclination, the corresponding polarizational charge with the density

$$-div\ P = -\left(\frac{\partial P_x}{\partial x} + \frac{\partial P_y}{\partial_y}\right)$$

is neutralized by the flowing volume charge with the density ρ. The free energy of the chiral smectic C* containing such a loop (or wall) is smaller, therefore this compensated wall must be stable and lose an ability to move, as in ordinary ferroelectrics.[5] As a consequence during the process of relaxation the appreciable part of electric charge is binded by disclinations and stabilizes the periodic system of walls. The another part of free charge carriers forms double-electric layers near the sample boundaries, inside these layers the field E_s exists. The essential electrostatic repulsion can exist between walls, it can exceed the repulsion conditioned by the orientational elasticity in strength.

The number of walls per unit length is of the order of h^{-1}. The free energy per unit area conditioned by the interaction between the subsurface field E_s and the helicoidal distribution of the polarization $P \sim \mu\theta$, less the energy of untwisted helix, is of the order of

$$PE_s\eta_s\left(\frac{r_D}{d} - \frac{\eta_s}{h_0}\right),$$

where $\eta_s \sim (K\theta/\mu E_s)^{1/2}$. The free energy density conditioned by the system of walls is

$$F_p \sim \frac{h_0}{h} P\sqrt{E_s E_c}\left(1 - \sqrt{\frac{E_s}{E_c}}\right)\left(\frac{r_D}{d}\right)^2. \tag{1}$$

The magnitude F_p is equal to zero at the critical value of untwisting field

$$E_c \sim \frac{K\theta}{\mu h_0^2}\left(\frac{d}{r_D}\right)^2.$$

The loss in energy conditioned by the mutual repulsion of the walls can be qualitatively written in the form

$$F_\rho = g\left(\frac{\rho}{h}\right)^2, \tag{2}$$

where the ratio (ρ/h) characterizes the gradient of charge density along the z-axis, g is the phenomenological parameter. The orientation elastic repulsion is of the order of

$$\frac{1}{h^2} K\theta^2 \exp\left(-\frac{h}{\eta_s}\right)$$

and it is assumed that it is small in comparison with the magnitude F_ρ.

The energy of disclinational loop in the polarizational helix, the plane of loop being parallel to the xy-plane and the axis of polarizational helix being parallel to z-axis, can be written in the form

$$F_d = bP^2 (R \ln(R/a) - qR^2),$$

where a is the radius of a disclination nucleus, b is the phenomenological parameter, R is the radius of loop. The magnitude F_d, analogous with the orientational energy of loop in a chiral smectic,[4] has the maximum $F^* \approx (b/4q)P^2 \ln(1/aq) \sim bhP^2$ corresponding to the radius value $R = R^* \sim q^{-1}$. The large loops are energetically favourable, but their appearance is determined by the probability to overcome the potential barrier F^* by thermal fluctuations, boundary conditions and various defects. One can write the number of disclinations per unit length along the axis of the helix (it is assumed that the disclinations overcome the barrier F^* by thermal activation) in the form

$$n = n_0 \exp(-F^*/T) = n_0 \exp(-h/\xi),$$

where n_0 is a parameter, $\xi \sim (T/bP^2)$. The free energy per unit length of large loop extended in the smectic plane (the length of the loop along the x-axis is L, the length of loop along the y-axis is equal to the thickness d is of the order of

$$-bP^2qd.$$

Corresponding the free energy per unit volume of the sample with the concentration of loops n can be written in the form

$$F_d \sim \frac{1}{h} bP^2 n. \tag{3}$$

The term (3) can exceed the corresponding term related to elastic orientational deformations for smectic with the large spontaneous polarization.

THE ELASTIC MODULATED STRUCTURE

Thus the qualitative estimation of the free energy density F for the sample of the chiral smectic C* with sufficiently large values of the spontaneous polarization, large concentrations of charge carriers and disclinations is described by the sum of Equations (1)–(3):

$$F = F_p + F + F_d = -\frac{h_0}{h} PE_s\tau + \frac{U^2}{2h^2} - \frac{\lambda}{h} P^2 \exp\left(-\frac{h}{\xi}\right), \tag{4}$$

where the dimensionless parameter

$$\tau = -\left(\frac{r_D}{d}\right)^2 \sqrt{\frac{E_c}{E_s}} \left(1 - \sqrt{\frac{E_c}{E_s}}\right)$$

characterizes a closeness of electric conditions to helix untwisting conditions, the parameter $\lambda \sim n_0 b$ has the dimensional representation of a length characterizing disclinations, the parameter $u = g^{1/2}\rho$ has the dimensional representation of the potential of electric field characterizing some electric conditions in vicinity of disclinations with the compensating charge. One should underline the qualitative character of the expression (4) where the subjects of interest are the dependence of the polarization P, the field E_s, the potential U (the charge density ρ) and the varied pitch h. The magnitude of h is found from the condition $\partial F/\partial h = 0$.

At active origin of loops, i.e., at $h < \xi$, one can see from (4) that the pitch is

$$h = \frac{U^2}{\lambda P^2 + h_0 P E_s^\tau} \tag{5}$$

If the field E_s is relatively weak, one can see from (1) that $PE_s^\tau \sim K\theta^2 h_0^{-2}$ and the inequality $h < \xi$ means that

$$\frac{U^2}{\lambda P^2 + K\theta^2 h_0^{-1}} < \frac{T}{bP^2},$$

i.e., $U^2 < Tn_0$. At these conditions there are the dependences of pitch and wave number on spontaneous polarization $h \sim P^{-2}$ and $q \sim P^2$ in the substances with large polarization.

If the activation process of the origin of loops is weak, i.e. at $h > \xi$ one derives from (4):

$$h = \frac{U^2}{h_0 PE_s^\tau + \lambda P^2 \exp(-U^2/\xi h_0 PE_s^\tau)} \tag{6}$$

At $U^2 < \xi h_0 PE_s^\tau$ Equation (6) transforms to Equation (5). The expressions (5) and (6) show the dependence of pitch on polarization is the complex function which is

conditioned by the electric potential of disclination (or the linear density of charge along the disclination) and the subsurface field E_s.
At the conditions

$$\xi h_0 P E_s^\tau >> U^2 >> h_0 \lambda P^2 >> h_0^2 P E_s^\tau,$$

which correspond to relatively large values of the correlation length ξ and the polarization P, the dependence

$$h \approx \frac{U^2}{\lambda P^2} \tag{7}$$

takes place, the high h being much larger than the equilibrium pitch h_0. At comparatively small polarization values

$$P << E_s \tau (h_0/\lambda)$$

one has the equation

$$h \sim \frac{U^2}{h_0 P E_s^\tau} \tag{8}$$

which preserves in increasing of the linear density of charge on the disclination, i.e., at $U^2 >> \xi h_0 P E_s^\tau$.

One can conclude from (5)–(8) that the structure period h substantially depends on the charge density (or the electric potential) on the disclination: $h \sim U^2$, $q \sim U^{-2}$. Consequently there is large ratio h/h_0 in samples with the high content of free carriers of charge, that was observed experimentally. The temperature dependences of pitch and wave number are determined by the temperature dependences of the polarization $P_0 = \mu\theta \sim \mu(T_c - T)^{1/2}$ and the equilibrium pitch $h_0(T)$. Near the transition point T_c the wave number $q \sim h^{-1}$, in accordance with (7), is

$$q \sim \frac{\mu^2 \lambda}{U^2} \left(1 - \frac{T}{T_c}\right) \tag{9}$$

If the field E_s is sufficiently small then, in accordance with (8) one has the equation

$$q \sim \frac{K}{U^2} q_0(T) \theta^2(T) \tag{10}$$

which results in more fast (in comparison with linear) decreasing of the number $q(T)$.

Figure 1 shows data on the temperature dependence of the wave number q for the static domain structure in two different mixtures, one can see a difference in the powers of corresponding dependences. For the mixtures 232 and 218 which

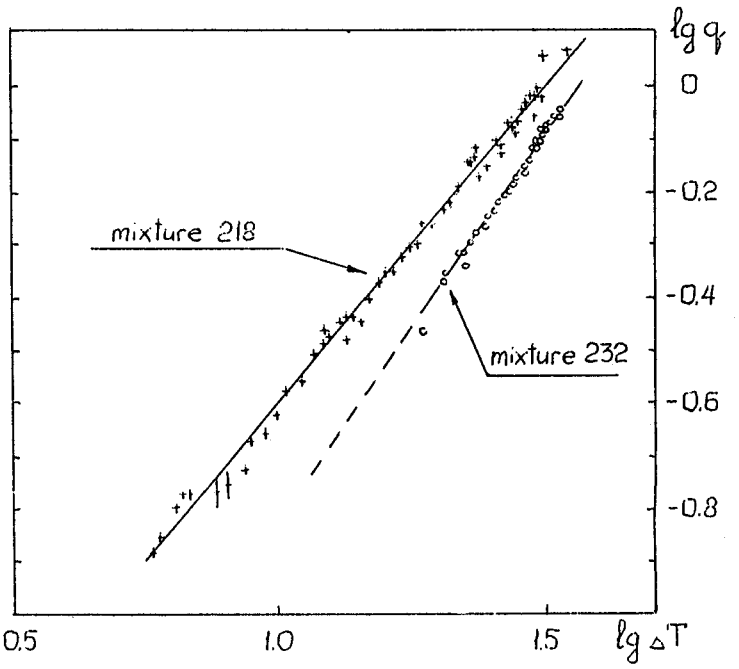

FIGURE 1 Temperature dependencies of the wave number of the static domain structure for mixtures 218 (+) (T_{AC} = 51.6°C, d = 29,5 mcm) and 232 (○) (T_{AC} = 51,8°C, d = 16,8 mcm).

have close values of the transition temperatures T_{AC} and of the spontaneous polarizations P, these powers appreciably differ, that is explained by the effect of chiral dopes. It is well known[6] that the last effect can result in increasing of the equilibrium helix pitch $h_0(T)$ which has the characteristic temperature dependence with a maximum near the phase transition point T_{AC}. In absence of untwisting conditions, i.e., at $E_s << E_c$ the ratio $(U^2/\xi h_0 P E_s^{\tau})$ in (6) is proportional to the magnitude $h_0(\mu U)^2$, the corresponding exponent factor in (6) rapidly decreases in increasing of the equilibrium pitch h_0, the piezocoefficient μ and the potential U. Therefore decreasing of chirality of mixtures must result in increasing the pitch h_0, fast decreasing of the exponent factor in (6) and in the temperature dependence (10) in which the magnitude $q_0(t)$ plays the noticeable role: it effectively increases the index in the power law $q(T)$.

It was noted in Reference 1 that the nonequilibrium pitch h is independent of the thickness d over a wide range of thicknesses (6–30 mcm), it is explained by the estimations (7)–(10). In the experiment the effect of increasing layer thickness manifests by an increase in the upper bound on a temperature up to which one can observe the nonequilibrium modulated structures. For example, in the mixture 218 with the phase transition point T_{AC} = 51,6°C one can oberve such a structure: at d = 6,4 mcm - up to T = 35°C, at d = 8,3 mcm - up to T = 38°C, at d = 29,5 mcm - up to T = 43°C. One can explain these data taking into account the dependence of E_c and τ on temperature and thickness [see Equation (4) and Equation (6)]. Really, an increase in the thickness d gives rise to fast increase in the threshold

value E_c and consequently at the fixed value E_s the condition of helix untwisting is executed at smaller values of the tilt angle θ, i.e., at temperatures more closely spaced to the threshold point T_{AC}. One can make an estimation of the dependence of the upper bound on a temperature T on the thickness d, taking into account the condition of untwisting $E_s = E_c$, where $E_c \sim \theta d^2$ and $\theta \sim (T_{AC} - T)^{1/2}$:

$$(T_{AC} - T_{\max}) \sim d^{-4}.$$

One should note that the $T_{\max}$ value rapidly brings close to T_{AC} value when the layer thickness increases. However in the experiment the upper bound on a temperature corresponds to possibilities of observation of such modulated structures: at $d > 30$ mcm the diffraction conditions worsen for structures with $h > d$, practically the upper bound is the temperature at which the period of structure is of the order of the layer thickness.

QUASISTATIONARY MODULATED STRUCTURE IN ALTERNATING ELECTRIC FIELD

The free charge, which compensates the polarization charge related to the unit length of disclination, can execute forced oscillations under the action of the alternating field $E = E_0 \exp(i\omega t)$. Consequently there are the macroscopic oscillations of polarization Δ = nuy near disclinations. One can qualitatively consider this system of charges on a disclination as a certain oscillator which is subjected to the action of the elastic form $-ky$ and the damping force $-m\gamma\dot{y}$, where m is an effective mass. The local field which excites such an oscillator includes two terms: $E + \beta\Delta$, where β is the Lorentzian factor. The equation of motion for the oscillator is[5]:

$$m\ddot{y} + 2m\gamma\dot{y} + ky = U(E + \beta\Delta),$$

or

$$\ddot{\Delta} + 2\gamma\dot{\Delta} + \omega_0^2\Delta = \frac{nU^2}{m} E_0 \cos(\omega t), \tag{11}$$

where $\omega_0^2 = (k - \beta nU^2)/m$.

One can estimate the ratio k/m in the equation for ω_0^2 supposing that charge carriers compensate polarizational charges in cylinders with the radius h and the length L, the disclination line being the cylinder axis. The density of these charges due to the flexoelectric effect[2] is of the order of $\rho \sim h^{-1}P \sim fh^{-2}$, where f is the flexoelectric coefficient. For small relative displacements of charges opposite signs under the action of field, there are the noncompensated charges with the density $\sigma \sim \rho y \sim fyh^{-2}$ on the cylinder surfaces. These charges change the energy of electric field in the cylinder on the average on the magnitude of the order of $\varepsilon^{-1}\sigma^2 V\varepsilon^{-1}\sigma^2 h^2 L \sim \varepsilon^{-1}f^2h^{-2}Ly^2$, where ε is the average dielectric permeability,

V is the cylinder volume. Thus the "elastic constant" k admits the estimation $k \sim \varepsilon^{-1} f^2 h^{-1} V$, so $k/m \sim \varepsilon^{-1} f^2 h^{-4} \tilde{\rho}$, where $\tilde{\rho} = m/V$ is the mass density. There is the proportionality $f \sim \theta$, i.e., $f^2 \sim \theta^2 \sim (T_c - T)$ in the chiral smectic C*.

The solutions of the Equation (11), for relatively small damping ($\omega_0^2 > 2\gamma^2$), have the resonance character at the frequency of alternating field $\omega = \omega_r = (\omega_0^2 - 2\gamma^2)^{1/2}$, the amplitude of oscillations being equal to

$$A_r = \frac{(nU^2/m)E_0}{2\gamma\sqrt{\omega_0^2 - \gamma^2}}$$

At the resonance conditions the average $\overline{P^2} = P_0^2 + \overline{\Delta^2}$ is equal to

$$\overline{P_r^2} = P_e^2 + \frac{(nU^2E_0/m)^2}{8\gamma^2(\omega_0^2 - \gamma^2)}. \tag{12}$$

The wave number q must be determined by the average $\overline{P^2}$ or $\overline{PE_s}$ in the equations (7) and (8). At the resonance conditions the magnitude q can appreciably increase. If $\omega_0^2 < 2\gamma$ then the amplitude A is maximum at $\omega = 0$ (static displacement) and decreases at frequency increasing:

$$(\overline{P^2} - P_0^2) \sim \begin{cases} (\gamma\omega)^{-1} & \text{at } \omega_0 << \omega << \delta \\ \omega^{-2} & \text{at } \omega >> \gamma \end{cases} \tag{13}$$

Thus one can see from (12) and (13) that at the resonance conditions the temperature dependence of q can remarkably differ from the proportionality $q(T) \sim P_0^2(T) \sim \theta^2(T)$. For the resonance increasing of the polarization inhomogeneity, which differs from the simple harmonic distribution, one can observe several orders of diffraction of light on such a structure. Figure 2 shows that the dynamic structure has the appreciably large wave number q near the phase transition point (in comparison with the static structure). This difference results in different values of powers in the power law $q(T_{AC} - T)$. It was noted that at temperatures larger than $T_{\max}$ the period of static structures becomes very large, the wave number of the dynamic structure must be determined by the second term in (12) and must be finite that is shown by Figure 2. At temperatures smaller than $T_{\max}$ the second term in (12) plays small role and correspondingly the wave numbers of static and dynamic structures are determined by the first term in (12), i.e., by the quadratic power in the spontaneous polarization that is also shown by Figure 2.

We observed also several orders of diffraction of light on such dynamic structures which existed in a narrow temperature interval that shows the existence of the mentioned resonance.

Generally speaking the existence of quasistationary spatially modulated structures under the action of alternating field is possible only at specific relationship between the field amplitude E_0 and the wave number q, such a relationship being dependent on specific physical mechanisms of phenomena.[2] The general feature of such phenomena is their threshold character, i.e., the finite values E_{oc} and q_c

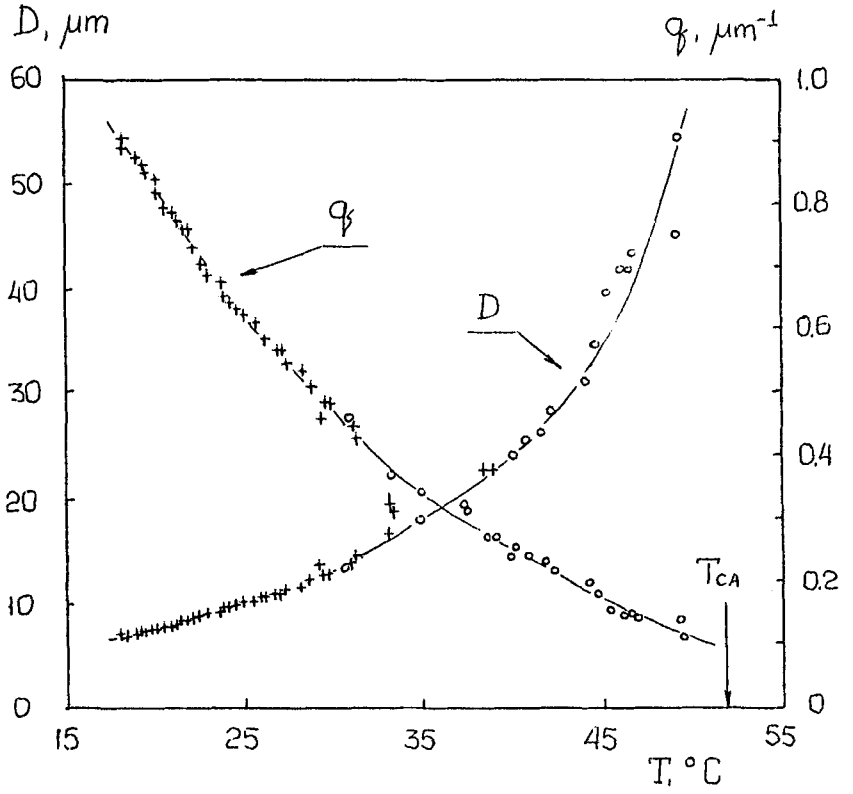

FIGURE 2 Temperature dependencies of the period h and the wave number q for the dynamic (○) and static (+) structures. Mixture 232, d = 16,8 mcm.

for the structure arising in the layer with rigid boundary conditions. For relatively soft boundary conditions (for example, for weak surface anchoring), one has the evident proportionality $E_0^2 \sim q^2$, i.e., the inhomogeneous structure does not arise in zero field. In considered case the periodic system of walls arises also in absence of field. Therefore, one can expect that the value of q conditioned by the mechanism discussed above can preserve if other physical reasons do not prevent this.

For substances with large polarization, the flexoelectric effect can be essential in the low-frequency field[2] if this effect is related to the modulation of the tilt angle $\theta(z)$ along the helix axis and to the inhomogeneity of the spatial distribution of solid charge $\rho(z)$. The corresponding term in the free energy density can be written in the form:

$$F_f = fE\rho \frac{\partial \theta}{\partial z}$$

Taking into account the terms F_f, $(1/2)K(\partial\theta/\partial z)^2$ and $(1/2)g(\partial\rho/\partial z)^2$ and varying the free energy by ρ and θ, one can obtain the equations

$$K\frac{\partial^2\theta}{\partial z^2} + fE\frac{\partial\rho}{\partial z} = 0$$
$$g\frac{\partial^2\rho}{\partial z^2} - fE\frac{\partial\theta}{\partial z} = 0 \qquad (14)$$

Supposing that $\theta \sim \rho \sim \exp(iqz)$, the magnitudes ρ and q being defined, one obtains from (14) that the nontrivial value of θ exists at the relationship

$$E^2 \sim (Kg/f^2)q^2 \qquad (15)$$

which determines the necessary value of field at the defined number q. The equation (15) shows that the amplitude of the low-frequency field, which does not destroy the structure by the flexoelectric effect, depends on temperature in such a way as the wave number q [see (9) and (10)].

Figure 3 shows the observed temperature dependence of the difference in potentials at which the dynamic inhomogeneous structure arises (at frequency 25 Hz). One can see that this dependence qualitatively repeats the dependence $q(T)$ taking into account the experimental error. Increase of the amplitude of alternating field

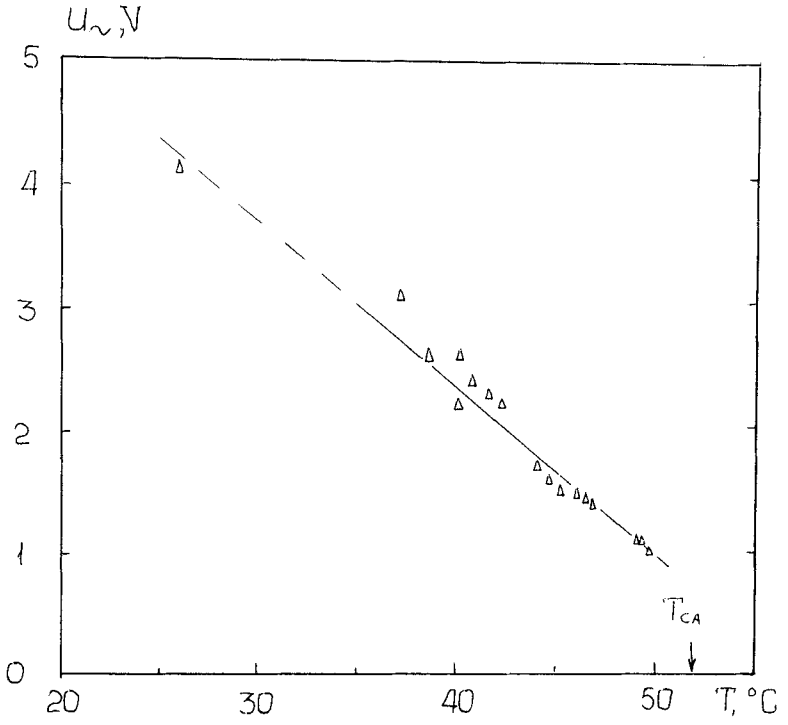

FIGURE 3 Temperature dependence of the electric potential U (frequency 25 Hz) at which the dynamic inhomogeneous structure arises. Mixture 232, d = 16,8 mcm.

above the threshold value results in appreciable decrease of the period of the modulated structure, i.e., in increase of the wave number q, that is explained by the proportionality (15). One should underline that the last relationship can take only at low frequencies of alternating field at which the flexoelectric effect is possible.[2] It was pointed above the resonance frequency in our experiments was sufficiently low.

CONCLUSIONS

The set of obtained data and made estimations shows that the observed phenomenon in liquid-crystalline ferroelectrics with sufficiently high spontaneous polarization is related to the formation of the nonequilibrium modulated polarizational and orientational structures in presence of the spatial electric charge with the modulated density. Relative stability of such structures is conditioned, probably, by screening of arising polarizational defects by free charge carriers. Low-frequency resonance character of these phenomena in alternating fields indicates the possible presence of the flexoelectric effect.

The temperature dependence of the frequencies ω_0 and ω_r is interesting, especially near the phase transition point, where these frequencies must decrease. Near T_{AC} point such oscillations must decay ($\omega_0^2 < 2\gamma^2$), therefore here the resonance does not take place. Far from the T_{AC} point one should make the estimation of ω_0 starting from the data on such macrocharacteristics as the flexoelectric coefficient $f \sim 10^{-5}$ CGSE units, the mass density $\tilde{\rho} \sim 1$ g.cm^{-3}, the period $h \sim 10^{-3}$ cm; $\omega_0 \sim fh^{-2}\,\tilde{\rho}^{-1/2} \sim 10$ sec^{-1}. For the temperatures close to the transition temperature, the frequency of own oscillations ω_0 has no sense because of large damping that is here probably the reason of the large experimental error in identification of the modulated structure.

The experimental estimation of the disclination electric potential U and of the corresponding density of free charge ρ is of great interest and demands special investigations.

References

1. L. A. Beresnev, M. V. Loseva, N. I. Chernova, S. G. Kononov, P. V. Adomenas and Y. P. Pozhydaev, *Pis'ma Zh. Eksp. Teor. Fiz.*, **51**, 457 (1990).
2. S. A. Pikin. *Structural transformations in Liquid Crystals,* Nauka, Moscow (1981) (in Russian).
3. M. Glogarova, J. Fousek, L. Lejcek and J. Pavel, *Ferroelectrics,* **58**, 161 (1984).
4. G. S. Chilaya, Z. M. Elashvili, M. A. Gogadze, S. P. Tavzarashvili, K. D. Vinokur and S. A. Pikin, *Liquid Crystals,* **5**, 1195 (1989).
5. D. Burfoot, *Polar dielectrics and their application.* Mir, Moscow (1981) (in Russian).
6. B. I. Ostrovskii, A. Z. Rabinovich, A. S. Sonin and B. A. Strukov, *Zh. Eksp. Teor. Fiz.*, **74**, 1748 (1978).