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HYDRODYNAMIC INSTABILITIES OF NEMATIC LIQUID CRYSTALS UNDER A. C. ELECTRIC FIELDS (*)

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Abstract. — We present an extension of the Helfrich theory of hydrodynamic instabilities to the case of alternating electric fields. The electrohydrodynamic effects are described by two coupled equations for the charge density \( q \) and the local curvature of the molecular alignment \( \psi \). The relaxation time for \( q \) is the dielectric relaxation time \( \tau \) (~ \( 10^{-2} \) s in typical samples). The relaxation time \( T \) for \( \psi \) is strongly dependent on the field magnitude. Provided that the sample thickness \( d \) is above a certain limit \( d_0 \), the nature of the instability is very different, depending on the ratio of the field frequency \( \omega \) to a critical frequency \( \omega_c \). For \( \omega < \omega_c \), the onset of instability corresponds to a charge \( q \) which oscillates at the frequency \( \omega_c \), while the curvature \( \psi \) is essentially time independent. For \( \omega > \omega_c \), the situation at threshold corresponds to a constant \( q \) and an oscillating \( \psi \). These predictions, together with the calculated curves of threshold voltage vs \( \omega \) and \( d \), are in reasonable agreement with a number of recent experiments.

I. Introduction. — The effects of electric fields on nematic liquid crystals are spectacular: when a nematic sample with negative dielectric anisotropy

\[ \varepsilon_{\parallel} - \varepsilon_{\perp} = \varepsilon_a < 0 \]

is placed between semi-transparent electrodes and a voltage \( V \) is applied, one observes the following sequence of events (1–4) (in increasing \( V_1 \)):

a) At a certain threshold \( V_e \) a cellular type of hydrodynamic motion appears: this is revealed by small concomitant deflexions of the molecular alignment, or by the motion of dust particles floating in the liquid.

b) At a higher threshold \( V_t \) the flow becomes turbulent; the molecular alignment is then strongly perturbed, and the scattering of visible light becomes very large: this type of effect has been called « dynamic scattering » by its discoverers.

In the present paper, we describe some theoretical calculations of the first threshold \( V_e \): this threshold is not the most interesting quantity from the point of view of optoelectronic applications, but it is the best probe available for fundamental studies. Many effects contribute to \( V_e \). A very lucid description of the state of affairs in given in the work of Rondelez [5]. Following his discussion, we must carefully distinguish the cases of DC and AC excitation:

1. For DC excitation. — The injection of charges via the electrodes plays an important role. This was already apparent in the work of the RCA group on \( V_e \). More recently, Koelmans and Van Boxtel [6] have shown that there remains a threshold \( V_{e1} \) (very similar in magnitude to \( V_e \) in the isotropic phase of the nematic: in this phase we are probably
dealing with the « Felici instability » of all liquid insulators under unipolar charge injection [7]. However, in the *nematic* phase, we must have some other processes superimposed to charge injection, since dynamic scattering is observed only for materials with $\epsilon_2 < 0$. The present interpretation of these processes is based on an idea of Carr [8] which has been worked out in great detail by Helfrich [9]; the principle is explained on figure 1. A qualitative discussion incorporating simultaneously the Felici effect and the Carr-Helfrich effect has been given recently [10].

![Figure 1](image.png)

**Fig. 1.** — A nematic liquid crystal is, at zero electric field, aligned along X-direction. A magnetic field $H$ and an AC external electric field $E(t)$ are respectively applied along $Y$ and $Z$-direction. $d$ is the sample thickness and $\varphi$ the fluctuation angle of the director (preferred direction) around X-direction.

2. AC Voltages. — At first sight this situation might appear more complex: in actual fact it is simpler, for the following reason: as soon as the frequency $\omega/2\pi$ exceeds a few cycles, charge injection becomes negligible [5]. Thus, by going to AC voltages, we eliminate all the complications related to the nature of the electrodes and the chemical reactions at their contact.

The present paper does not take into account any injection effects: it is thus restricted to the AC regimes. The only new contribution which we bring in, as compared to the work of Helfrich [9], is to introduce time dependent equations, where the time lags associated with the charge response and with theorientational response of the liquid crystal are taken into account. Just as in ref. [9] our description contains certain oversimplifications: if the sample is parallel to the $xy$ plane, with the molecules along $x$ in the unperturbed state and the electric field along $z$ (see Fig. 1), we assume that all physical quantities depend only of $x$, not of $y$ or $z$: the only type of deformation which is taken into account is a pure bending mode. This is qualitatively correct, as shown by a number of recent experiments. However, near the walls, the «director» $\mathbf{n}$ must also depend on $z$ to satisfy the correct boundary conditions. In our approximate treatment, this is not taken into account; we simply say, as in ref. [9], that the wavelength of the fluctuating mode along $x$ cannot be much larger than the sample thickness ($^1$). A more detailed three dimensional analysis could be attempted, but the present, rough approach allows for much more physical insight.

In section II, we rederive the fundamental dynamic equations describing the charge flow and the changes in molecular orientation, based on the hydrodynamic theories of Ericksen [11] and Leslie [12], [13]. (We consider only materials where the anisotropy of the electrical conductivity $\sigma_2 = \sigma_0 - \sigma_1$ is positive — as it has always been found up to now.) In section III, we apply these equations to the low frequency regime, where the changes are modulated at the frequency $\omega$ of the electric field. We call this the « conducting regime ». In section IV, we discuss the high frequency regime where the charge distribution becomes essentially static (« non conducting regime »). In section V we discuss some extension of our results to systems where the dielectric anisotropy would be slightly positive, rather than negative, but where the unperturbed molecular alignment is still in the plane of the field (because of the walls or because of a magnetic field $H$ applied along $x$). The case where the applied electric field, instead of being a sine wave, is a succession of rectangular pulses, will be the object of a second paper: it leads to a very different diagram of threshold versus frequency, and could be of some experimental interest.

II. Electrohydrodynamic equations. — As explained in the introduction, we consider a nematic slab of thickness $d$, lying in the $xy$ plane (Fig. 1) and submitted to an electric field $E = E_m \cos \omega t$ along $z$. We assume that in the unperturbed state the molecules are aligned along the $x$ direction, and we also allow for a stabilising magnetic field $H$ applied along $x$. We consider pure bending fluctuations ($^2$) where the director $\mathbf{n}$ is in the $(x, z)$ plane, making a small angle $\varphi$ with the $x$ direction. Finally we assume that $\varphi$ is only a function of $x$.

Associated with such an orientational fluctuation, we have a space charge $q(x)$ and a hydrodynamic flow velocity: in our approximation everywhere things depends only on $x$, this velocity has only one non-zero component $v_x(x)$ ($^3$).

For the internal rotation, we can neglect the inertial terms. Let us define as in ref. [14] the viscous torque as the frictional torque exerted by the molecules on the over-all hydrodynamic motion. This viscous torque must be equal to the sum of the elastic,

(1) For some cases, we will find that the wavelength is in fact much smaller than the thickness: then our analysis is probably quite correct.

(2) Helfrich has shown that, for this orientation of the electric field, these pure bending fluctuations are the only ones to induce shear flows.

(3) For the slow motions of interest, the nematic may be considered as incompressible: here this imposes $\partial v_x/\partial x = 0$ and finally $v_z = 0$ when there is no general drift along $x$. 

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**Footnotes:**

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**Figure 1:** A nematic liquid crystal is, at zero electric field, aligned along X-direction. A magnetic field $H$ and an AC external electric field $E(t)$ are respectively applied along Y and Z-direction. $d$ is the sample thickness and $\varphi$ the fluctuation angle of the director (preferred direction) around X-direction.
magnetic and dielectric torques exerted on the molecules:

$$\Gamma_{\text{magn.}} = \Gamma_{\text{el.}} + \Gamma_{\text{magn.}} + \Gamma_{\text{diel.}} . $$

The dielectric torque is induced by the total electric field, i.e. the sum of the external field and the field due to space charges $q$.

**II.1 CONTRIBUTIONS TO THE TORQUE.** — The electric, dielectric and magnetic torques can be derived from the free-energy

$$F = \int \int g(r) dv$$

$$g(r) = \frac{1}{2} \left\{ K_{11} (\text{div} \ n)^2 + K_{22} (\n \cdot \nabla \phi (n))^2 + K_{33} (\n \times \nabla \phi (n))^2 - \frac{\varepsilon_0}{4\pi} (n \cdot E)^2 - \chi_0 (n \cdot H)^2 \right\}$$

where $K_{ii}$ are the elastic coefficients in Franck's [15] notation, $\chi_0$ and $\varepsilon_0$ the anisotropic parts of the magnetic susceptibility and dielectric constants

$$\chi_0 = \chi_0 - \chi_\perp ; \quad \varepsilon_0 = \varepsilon_0 - \varepsilon_\perp .$$

It will be assumed that the fluctuation angle, $\phi$, is small: only first order terms in $\phi$ will be kept. This fluctuation induces a space charge $q$ which in turn induces the $x$-component of $E_0$, $E_x$ of the total electric field [9]. As shown by Helfrich, $E_x$ and $q$ are first-order terms in $\phi$. For the sake of simplicity, we will omit the part of the $z$-component of the electric field induced by the space-charge $q$; it can be easily shown that this component occurs only in second-order terms in $\phi$. Hence $E_z = E(t)$.

In our frame of reference the torques are given by the functional derivative

$$\Gamma_r (r) = \frac{\delta F}{\delta \phi(r)} .$$

Assuming a periodic fluctuation

$$\phi = \phi_0 \cos kx$$

one finds, for the elastic, magnetic and dielectric torques:

$$\Gamma_{\text{el.}} = -K_{33} \frac{\delta^2 \phi}{\delta x^2} = K_{33} k^2 \phi$$

$$\Gamma_{\text{magn.}} = \chi_0 H^2 \phi$$

$$\Gamma_{\text{diel.}} = -\frac{\varepsilon_0}{4\pi} E^2 \left( \phi + \frac{E_x}{E} \right) .$$

Using the notations of reference [14], the viscous torque is

$$\Gamma_{\text{visc.}} = n \times \left( \gamma_1 n + \gamma_2 A \cdot n \right)$$

with

$$N = \left( \Omega - \frac{1}{2} \nabla \phi (n) \right) \times n$$

$$A_{ij} = \frac{1}{2} \left[ \frac{\partial \psi_i}{\partial r_j} + \frac{\partial \psi_j}{\partial r_i} \right]$$

where the director $n(r)$ is a unit vector along the preferred orientation of the molecule; $\Omega$ is the angular velocity of the director and $v$ is the fluid velocity. Recalling that the only non-zero component of $\psi$ is $\varepsilon_0 (x, t)$ the only non-zero components of tensor $A$ are

$$A_{xx} = A_{xx} = \frac{1}{2} \frac{\partial \psi_x}{\partial x} .$$

Hence

$$\Gamma_{\text{visc.},y} = -\gamma_1 \left[ \phi - \gamma_1 - \gamma_2 \frac{\partial \psi_x}{\partial x} \right] .$$

(II, 3)

We will use further the following notations

$$D_0^2 = -\varepsilon_0 E_0^2 = 4\pi \frac{\varepsilon_0}{\varepsilon_1} \left[ \chi_0 H^2 + K_{33} k^2 \right] .$$

(II, 4)

The minus sign in the definition of $E_0^2$ will be useful in sections III and IV where we assume a negative dielectric anisotropy, which is the case for the experimental data. In section V, we will allow $\varepsilon_0$ to be positive, and use only $D_0^2$.

Inserting eqs. (II, 2-4) into eq. (II, 1), one finds

$$-\Gamma_{\text{visc.},y} = \gamma_1 \left[ \phi - \gamma_1 - \gamma_2 \frac{\partial \psi_x}{\partial x} \right]$$

$$= \frac{\varepsilon_0}{4\pi} \left[ \left( \frac{\varepsilon_1}{\varepsilon_0} E^2 + E^2 \right) \phi + EE_x \right] .$$

(II, 5)

**II.2 CHARGE-BALANCE.** — The conservation equation for charges is

$$q + \frac{\partial J_x}{\partial x} = 0$$

(II, 6)

where $q$ is the excess charge per unit volume and $J$ the electric current. It is shown in Appendix A that the diffusion currents can be neglected. Then

$$J = \sigma_1 E + \sigma_2 n (n \cdot E) .$$

(II, 7)

Keeping only first-order terms in $\phi$ and $E_x$,

$$J_x = \sigma_1 E_x + \sigma_2 \varepsilon_0 E \phi .$$

(II, 8)

From the relation

$$\frac{\partial D_x}{\partial x} = \frac{\partial \psi}{\partial x} (\varepsilon_0 E_x + \varepsilon_0 E \phi) = 4\pi q$$

(II, 9)

one obtains

$$\frac{\partial E_x}{\partial x} = \frac{4\pi q}{\varepsilon_0} = \frac{\varepsilon_0}{\varepsilon_0} E \psi$$

(II, 10)

where $\psi = \partial q/\partial x$ is the local curvature. Inserting eqs. (II, 8) and (II, 10) into eq. (II, 6), one has the charge-balance equation

$$\dot{\psi} + \frac{\dot{q}}{\varepsilon_0} + \sigma_1 E \psi = 0 .$$

(II, 11)
Here \( \tau \) is the dielectric relaxation time for charges. Eq. (II, 11) relates the charge \( q \) to the curvature \( \psi \).

II.3 ACCELERATION EQUATION. — The equation of motion is (4)

\[
\rho \frac{\partial \tilde{v}}{\partial t} = - \text{Div} (\rho \tilde{v}) + \text{Div} (\sigma - \bar{\sigma}) + qE \quad (II, 13)
\]

where \( \bar{\sigma} \) and \( \tilde{\sigma} \) are the elastic and viscous stress tensor \([11], [15]\).

\[
\sigma_{ij} = - p \delta_{ij} - \sum_{i} \left( \frac{\partial g_{it}}{\partial n_{i}} \right) n_{j,i} \]
\[
g_{it} = \frac{1}{2} \left[ K_{11}(\text{div} n)^2 + K_{23}(n \cdot \text{curl} n)^2 \right] \quad (II, 14)
\]

We use for \( \tilde{\sigma} \) the notation of ref. [16]: this takes into account the effect of the Onsager reciprocal relations on the expression of the viscous stress tensor. \( N \) and \( \bar{A} \) are defined in section (II, 1).

The \( x \)-component of eq. (II, 13) defines the pressure \( p \). We are interested only in the \( z \)-component of this equation. Since \( v_{x} = 0 \) and \( n_{i,z} = 0 \),

\[
(D \text{Div} \tilde{v}) = \frac{\partial}{\partial x} (\rho v_{x} v_{z}) = 0
\]

\[
(D \text{Div} \sigma_{zz}) = \frac{\partial}{\partial x} \sigma_{zz} = \frac{\partial}{\partial x} \sum_{i} \left( \frac{\partial g_{it}}{\partial n_{i}} \right) n_{0,i} = 0.
\]

Hence

\[
\rho \frac{\partial v_{z}}{\partial t} = \frac{\partial}{\partial x} \sigma_{zz} + qE. \quad (II, 15)
\]

Keeping only first-order terms in eq. (II, 14), \( \sigma_{zz} \) is given by

\[
\sigma_{zz} = + \left( \frac{\gamma_{1} - \gamma_{2}}{2 \gamma_{1}} \right) \Gamma_{\text{visc.}} + \eta' \frac{\partial v_{z}}{\partial x}. \quad (II, 16)
\]

where

\[
\eta' = \frac{\gamma_{1} (2 \alpha^{2} + \beta)}{4 \gamma_{1}}. \quad (II, 17)
\]

From thermodynamical relations [16], \( \eta' \) is positive definite. Inserting (II, 16) in (II, 15), we find, for the acceleration equation:

\[
\rho \frac{\partial v_{z}}{\partial t} = \frac{\gamma_{1} - \gamma_{2}}{2 \gamma_{1}} \frac{\partial}{\partial x} \Gamma_{\text{visc.}} + \eta' \frac{\partial^{2} v_{z}}{\partial x^{2}} + qE. \quad (II, 18)
\]

In this equation the left h. s. may be neglected (a) in the « conduction » regime where this term vanishes (b) in the « dielectric » regime where \( \rho \partial v_{z} / \partial t \) is of order \( \alpha \nu \nu_{z} \) and can be neglected with respect to

\[
\eta' \frac{\partial^{2} v_{z}}{\partial x^{2}},
\]

which is of order \( \eta' k^{2} \nu_{z} \) (Typical experimental conditions are \( \omega \approx 10^{-4} - 10^{4} H \); \( k \approx 2 \times 10^{4} \text{cm}^{-1} \); \( \eta' \approx 10^{-1} \text{poise} \)).

II.4 EQUATION FOR THE CURVATURE. — Using eqs. (II, 2), (II, 3), (II, 10) and (II, 18), the \( x \)-derivatives of the torques can now be expressed as functions of the charge \( q \) and the curvature \( \psi \). Let us introduce the following notations:

\[
\eta_{0} = \frac{\left( 2 \gamma_{1} \right)^{2}}{\gamma_{1} - \gamma_{2}} \eta'.
\]

\[
\frac{\partial}{\partial x} \Gamma_{\text{visc.}} = \frac{\gamma_{1} \eta_{0}}{\gamma_{1} + \eta_{0}} \left[ \psi + \frac{1}{\eta_{0}} \frac{2 \gamma_{1}}{\gamma_{1} - \gamma_{2}} qE \right]. \quad (II, 21)
\]

Using eqs. (II, 2), (II, 4) and (II, 10), one finds

\[
\frac{\partial}{\partial x} \Gamma_{\text{visc.}} = \frac{\gamma_{1} \eta_{0}}{\gamma_{1} + \eta_{0}} \left[ \psi + \frac{1}{\eta_{0}} \frac{2 \gamma_{1}}{\gamma_{1} - \gamma_{2}} qE \right]. \quad (II, 21)
\]

Using eqs. (II, 2), (II, 4) and (II, 10), one has

\[
\frac{\partial}{\partial x} \Gamma_{\text{visc.}} = \frac{\gamma_{1} \eta_{0}}{\gamma_{1} + \eta_{0}} \left[ \psi + \frac{1}{\eta_{0}} \frac{2 \gamma_{1}}{\gamma_{1} - \gamma_{2}} qE \right]. \quad (II, 21)
\]

Hence, from eq. (II, 1) one finds the curvature equation

\[
\frac{\partial}{\partial x} = \frac{\gamma_{1}}{\gamma_{1} + \eta_{0}} \left[ \frac{\psi}{\gamma_{1} + \eta_{0}} + \frac{\lambda E^{2}}{\eta} \psi + \frac{qE}{\eta} \right] = 0 \quad (II, 24)
\]
is an effective viscosity. From theoretical considerations on flow-alignment of nematic liquid crystals, one can deduce the following inequalities \[13\]

\[\gamma_1 - \gamma_2 > 2 \gamma_1 > 0 ; \quad 0 < - (\gamma_1 + \gamma_2) < 1 .\]

Hence \(2 \gamma_1 / (\gamma_1 - \gamma_2)\) is of order unity. For small values of \(\varepsilon \approx (\varepsilon_0, \varepsilon_0') \approx 1\),

\[\eta \approx \eta_0 \approx \eta' \approx 2 \varepsilon_0 + \beta - \gamma_1 \frac{4}{\eta_0} .\]

On the contrary, for materials having a strong positive dielectric anisotropy, \(\eta\) can be negative. \(1/\eta\) vanishes for

\[\varepsilon_0 \approx \varepsilon_0' \approx \frac{2 \varepsilon_0 + \beta - \gamma_1}{4} .\]

We shall see in section V the consequences of this point.

Eq. (II, 24) can now be rewritten as:

\[\psi + \frac{1}{T} \psi + \frac{q E}{\eta} = 0 \quad (\text{II, 27})\]

where the time-dependent quantity

\[\frac{1}{T} = \lambda(E^2 + E^2_0) \quad (\text{II, 28})\]

appears as a relaxation rate for the curvature.

11.5 FORMAL MATHEMATICAL PROPERTIES. — Eq. (II, 11) and (II, 27) are two coupled equations for the charge density \(q\) and the curvature \(\psi\). They are linear in \(q\) and \(\psi\); from a mathematical point of view, they have some similarities with the spinor equations describing the motion of a quantum mechanical spin \(S = \frac{1}{2}\) in a radiofrequency field. However, one further complication is present: the relaxation rate \(1/T\) for the curvature is not a constant, because \(E\) is modulated in time. Let assume a sinusoidal field \(E = E_M \cos \omega t\) and rewrite our two equations for this case

\[
\dot{q} + \frac{q(t)}{\tau} + \sigma_H E_M \psi(t) \cos \omega t = 0 \quad (\text{II, 29})
\]

\[
\dot{\psi} + \lambda(E^2_0 + E^2_M \cos^2 \omega t) \psi(t) + \frac{E_M}{\eta} q(t) \cos \omega t = 0 . \quad (\text{II, 30})
\]

We shall also find the following notations useful

\[
\begin{align*}
\frac{1}{T_f} &= \frac{\lambda E^2_M}{2} = \lambda E^2(t) \\
\mu^2 &= \frac{E^2_0}{E^2_M} . \quad (\text{II, 31})
\end{align*}
\]

The system (II, 29), (II, 30) accepts solutions of the form

\[
\begin{bmatrix}
q(t) \\
\psi(t)
\end{bmatrix} = \begin{bmatrix}
a(t) \\
b(t)
\end{bmatrix} e^{i\nu} \quad (\text{II, 32})
\]

where \(a(t)\) and \(b(t)\) are periodic functions of time with period \(2 \pi/\omega\), and \(s\) is a number which may in principle, be real or complex. The proof of eq. (II, 31) is similar to the proof of the Bloch-Floquet theorem for electron propagation in one dimensional solids and is given in appendix B. For each value of \(E_M\), there are two independent solutions of the form (II, 32) with different parameters \(s_1(E_M)\) and \(s_2(E_M)\). We shall now prove that when the inequalities (III, 1) are satisfied, the curvature \(\psi(t)\) at threshold is essentially independent of time. Our starting point is eq. (II, 30) which may be rewritten as

\[
T_f \dot{\psi} + (2 \mu^2 + 1 + \cos 2 \omega t) \psi(t) = -\frac{2}{\lambda \eta E_M} q(t) \cos \omega t . \quad (\text{III, 2})
\]

This may be transformed into the integral form
\[ \psi(t) = -\frac{2}{\lambda\eta E_M T_f} \exp \left[ -\left(1 + 2 \mu^2\right) \frac{t}{T_f} - \frac{\sin 2 \omega t}{2 \omega t} \right] \times \int_{-\infty}^{t} dt' q(t') \cos \omega t' \times \exp \left[ \left(1 + 2 \mu^2\right) \frac{t'}{T_f} + \frac{\sin 2 \omega t'}{2 \omega t'} \right] \]

\[ \text{III.3 DISCUSSION.} \quad \text{The threshold field depends only on two parameters which characterize the nematic material. The first one is the relaxation time for charges, } \tau. \text{ The second one is the dimensionless coefficient } \xi. \]

\[ \xi^2 = \frac{\sigma_H \tau}{\lambda\eta}. \quad \text{(III.8)} \]

Inserting eq. (II, 31) into eq. (III, 7), one finds for the threshold field

\[ \frac{E^2(t)}{2} = \frac{E_{\text{DC}}^2(1 + \omega^2 \tau^2)}{\xi^2 - (1 + \omega^2 \tau^2)}. \quad \text{(III.9)} \]

III.3 DISCUSSION. — The threshold field depends only on two parameters which characterize the nematic material. The first one is the relaxation time for charges, \( \tau \). The second one is the dimensionless coefficient \( \xi^2 \).

\[ \xi^2 = \left[1 - \frac{\sigma_1 \sigma_2}{\sigma_1 \sigma_2}ight] \left[1 - \frac{\epsilon_1 \gamma_1}{\epsilon_1 (\gamma_1 - \gamma_2)(\gamma_1 + \eta_0)} \right] \quad \text{(III.10)} \]

For M B B A \( (\sigma_1/\sigma_2 = 1.5; \epsilon_1 = 4.7; \epsilon_2 = 5.4; \eta_0/\gamma_1 \approx 0.4) \) one finds \( \xi^2 \approx 3.2 \). The role of the parameter

\[ \theta_H = \epsilon_1(\xi^2 - 1) \]

has been pointed out by Helfrich [9] for the DC regime. For this reason we shall refer to \( \theta_H \) as the Helfrich parameter. It will play a special role in the discussion on the instabilities for negative or positive \( \epsilon_1 \) which is given in section V.

It is interesting to note that (in spite of the underlying approximation \( \omega T_f \gg 1 \)) eq. (III, 9) gives the exact threshold in the limit \( \omega \to 0 \). This DC threshold can be directly obtained from eq. (II, 29-30):

\[ E_{\text{DC}} = \frac{E_0^2}{(\xi^2 - 1)}. \quad \text{(III.11)} \]

For increasing \( \omega \), eq. (III, 9) leads to a « cut-off frequency »

\[ \omega_c = \frac{(\xi^2 - 1)^{\frac{1}{2}}}{\tau}. \quad \text{(III.12)} \]

For \( \omega \) increasing toward \( \omega_c \), the threshold field tends to become very large; in fact eq. (III, 9) becomes incorrect when \( E_M \) reaches a value \( E_{\text{Me}} \) such as \( 1/T_f \) becomes comparable to \( \omega \). On figure 2 is shown the range of validity of eq. (III, 9). The approximation will be valid up to \( \omega = \omega' \), \( \omega' \) will be close to \( \omega_c \) if \( \lambda E_0^2/\omega_c < 1 \). Let us discuss this quantitatively for \( H = 0 \) and assuming values of \( (\xi^2 - 1) \) of order unity. Then, if \( d \) is the sample thickness, replacing (tentatively) the wave-vector \( k \) by \( n/d \) as in reference [9], we have

\[ \frac{\lambda E_0^2}{\omega_c} \approx \frac{\pi^2 K_{13} \tau}{\eta^2} < 1 \]

where

\[ \frac{1}{\eta} = \frac{1}{\gamma_1} + \frac{1}{\eta_0}. \]
FIG. 2. — The full line gives the threshold-field of instability versus $\alpha^2$ in reduced coordinates. The dashed line is the parabola $u = (\zeta^2 - 1) (d/d)^4 v^2$ corresponding to the condition $\omega T_\gamma = 1$. The two lines intersect at points $(u_0, v_0)$ and $(u_\epsilon, v_\epsilon)$. The assumption $\omega T_\gamma > 1$ is valid in the corresponding frequency range $(\omega_0, \omega_\epsilon)$.

Thus one condition for the existence of the low frequency, conducting regime is

$$d > d_c$$

where $d_c = \pi(K_{33} \tau/\eta)^{1/2}$. Typically, $K_{33} \approx 10^{-6}$ dynes, $1/\tau \approx 2 \times 10^2$ s$^{-1}$, $\eta' = 10^{-1}$ poise; this gives $d_c \approx 7 \mu m$.

It must be noted that, as can be seen on figure 2, this condition ensures the validity of eq. (III, 9) not only up to frequency $\omega = \omega_\epsilon$ but also in the low frequency range. For $d/d_c \ll 1$, one finds for the limiting frequencies

$$\omega_0 = \frac{\omega_\epsilon}{(\zeta^2 - 1)^{1/2}} \left( \frac{d_c}{d} \right)^2;$$

$$E_{0}^2 = \frac{E_{\epsilon}^2}{\zeta^2 - 1} \left[ 1 + \frac{\zeta^2}{(\zeta^2 - 1)^3} \left( \frac{d_c}{d} \right)^4 \right];$$

$$\omega_\epsilon' = \omega_\epsilon \left[ 1 - \frac{\zeta^2}{2(\zeta^2 - 1)^{3/2}} \left( \frac{d_c}{d} \right)^2 \right];$$

$$E_{0E}^2 = \frac{E_{0}^2}{(\zeta^2 - 1)^{1/2}} \left( \frac{d_c}{d} \right)^2.$$  

Taking $\zeta^2 = 3$ ; $1/\tau = 200$ s$^{-1}$ ; $d = 100 \mu m$ ; $d_c = 7 \mu m$, one finds $\omega_0/2 \pi \approx 0.1$ Hz ; $\omega_\epsilon/2 \pi \approx 50$ Hz ; $$(\omega_\epsilon - \omega_\epsilon') \approx 0.5$ Hz .

Thus, for materials having a negative dielectric anisotropy, the « conduction regime » occurs in the frequency range $(0, \omega_\epsilon)$. (For very low frequencies further complications might be caused by carrier injection, but we do not consider this here.)

We shall find in section IV that, for frequencies $\omega > \omega_\epsilon$, another regime occurs, the so-called « dielectric » or « fast turn-off » regime. $\omega_\epsilon$ thus appears as a limiting frequency between two regimes rather than a real cut-off frequency.

We shall now discuss briefly the dependence of the threshold on thickness and on magnetic field. The equivalent field $E_0$ is given by eq. (II, 4)

$$E_0^2 = -\frac{4}{e_a e_{11}} \left( \chi_a H^2 + K_{33} \pi^2 \right)$$

where $k = k_x$ is the wave-vector along $x$-axis. Both experimental conditions and qualitative physical arguments suggest that $k \approx \pi/d$ where $d$ is the sample thickness. Hence

$$E_0^2 = -\frac{4}{e_a e_{11}} \frac{K_{33} \pi^2}{d^2} \left[ 1 + \frac{d^2}{\pi^2 \zeta^2} \right].$$

a) For low magnetic fields ($\zeta > d$). — From eqs. (III, 9) and (III, 14), one finds, for the rms threshold voltage $V_{thr}$,

$$V_{thr}^2(\omega) = \frac{V_0^2(1 + \omega^2 \tau^2)}{\zeta^2 - 1 - \omega^2 \tau^2}$$

where

$$V_0 = \left[ -\frac{4}{e_a e_{11}} K_{33} \pi^2 \right]^{1/2}.$$

Eq. (III, 15) shows that, for low magnetic fields, the experiments will give a voltage threshold independent of sample thickness. This is well confirmed by the experimental results [1], [2].

b) For higher magnetic fields. — One has

$$V_{thr}(\omega, H) = V_{thr}(\omega, 0) \sqrt{1 + \frac{d^2}{\pi^2 \zeta^2}}$$

or

$$V_{thr}(\omega, H) = V_{thr}(\omega, 0) \sqrt{1 + \frac{H^2}{H_0^2}}$$

where

$$H_0 = \frac{\pi}{d} \sqrt{\frac{K_{33}}{\chi_a}}$$

(with $d = 100 \mu m$, $H_0 \approx 10^3$ G).

For $H > H_0$, $V_{thr}(\omega, H)$ will be proportional to $d$ ; one will observe, for fixed $\omega$, a field threshold proportional to $H$. These results are in reasonable agreement with the experiments made at Orsay [5].
IV. « Dielectric régime » (ωτ ≫ 1). — When ωτ is large, the charges accumulated by the Carr-Helfrich process do not have enough time to flow during one cycle: the charge density q (at threshold) becomes time independent. This may be proved explicitly on an integral version of eq. (II, 29):

\[ q = - \sigma_H E_M e^{-t'/\tau} \int_{-\infty}^{t'} e^{t'/\tau} \psi(t') \cos \omega t' \, dt'. \]  

(IV, 1)

In the limit ωτ ≫ 1, \( \exp(t'/\tau) \) can be taken as a constant over one period. Thus we can replace in the integral \( \psi(t') \cos \omega t' \) by its average value over a period. Expanding \( \psi(t) \) as a Fourier series

\[ \psi(t) = \sum_{n=0}^{\infty} \psi_n \cos n\omega t + \psi_n \sin n\omega t \]

one has

\[ <\psi(t) \cos \omega t>_{\text{per.}} = \frac{\psi_1}{2} \]

and

\[ q = - \frac{\sigma_H E_M \tau}{2} \psi_1. \]

The space-charge is time independent. This regime is a « non conducting » regime. Using eq. (IV, 2), eq. (II, 30) is easily integrated:

\[ g(X, Y) \]

is a continuous function of X and Y. Let us rewrite eq. (II, 30) and (IV, 4) as:

\[ \begin{align*}
\psi(u) + 2 \psi(u) (Y + X \cos^2 u) - \zeta^2 X \psi'_1 \cos u &= 0, \\
\psi'_1 &= \frac{1}{\pi} \int_0^{2\pi} \psi(u) \cos u \, du.
\end{align*} \]

(IV, 9) (IV, 10)

where \( u = \omega t \) and \( \psi'(u) = \delta \psi/\delta u \).

From these last equations it can be easily shown that

(i) When \( X \to 0 \), the condition for \( \psi \) to be periodic is

\[ \lim_{X \to 0} \frac{1}{\zeta^2} = \frac{\lim g(X, Y)}{X \to 0} = \frac{Y \mathcal{X}}{4Y^2 + 1} (Y \neq 0) \]

Thus, for \( X \to 0 \), \( g(X, Y) \to 0 \).

(ii) When \( X \to \infty \), from eq. (IV, 9), for finite Y, \( \psi(u) \to \psi(\zeta^2/2 \cos u) \) and from eq. (IV, 10)

\[ \frac{1}{\zeta^2} = \lim_{X \to \infty} g(X, Y) = 1. \]

(iii) If X or Y is negative, the integral in eq. (IV, 3) diverges.

Hence for any value of \( \zeta^2 (\zeta^2 > 1) \) (5) one can find solutions of the equation

\[ g(X, Y) = \frac{1}{\zeta^2} \]

(IV, 11)

and the curve defined by this equation is located in the first quadrant \( (X, Y > 0) \).

Let \( X_m \) and \( Y_m \) be the values of X and Y that give the absolute minimum of X on this curve. \( X_m \) and \( Y_m \) correspond to the threshold conditions and depend only on \( \zeta^2 \). This leads to two requirements, listed below under (a) and (b):

(a) Threshold field.

\[ \frac{E_M^2}{2} = \frac{E^2(t)}{\lambda} = \frac{X_m(\zeta^2)}{\lambda} \omega. \]

(IV, 12)

From eq. (IV, 12), \( \overline{E^2(t)} \) appears to be proportional to \( \omega \). This is an asymptotic result for high \( \omega \). At slightly lower frequencies, the correct result is \( E^2 = \text{constant} + \omega X_m/\lambda \). We shall now discuss briefly an improved approximation leading to this form.

We must take into account the fact that \( X_m \) depends not only on \( \zeta^2 \) but also on \( \omega \). For \( \omega \tau \gg 1 \), we can expand \( X_m \) on powers of 1/\( \omega \tau \):

\[ X_m(\zeta^2, \omega \tau) = X_m(\zeta^2) + \frac{1}{\omega \tau} X'_m(\zeta^2) + O(\omega^2 \tau^2). \]

(5) It appears, from computer calculations, that, eq. (IV, 11) has no solution for \( \zeta^2 < 1 \). We have not been able to give a mathematical demonstration of this point. It implies that there is no « dielectric » instability for \( (\zeta^2 - 1) < 0 \).
Then we have

$$\bar{E}^2(t) = \frac{1}{\lambda} \left[ \omega \chi_m^0(\zeta^2) + \frac{1}{\tau} \chi_m^0(\zeta^2) \right]$$  \hspace{1cm} (IV, 13)

and we find a straight line for the plot of $\bar{E}^2$ versus $\omega$. This is in good agreement with the experimental result [5], [17], [18]. Finally, inserting eq. (II, 19) in eq. (IV, 13) one obtains

$$\bar{E}^2(t) = \left[ \frac{\varepsilon_\parallel}{\varepsilon_\perp} \frac{4 \pi \eta_1 \eta_0}{(\tau_1 + \tau_0)} \chi_m^0(\zeta^2) \right] \times$$

$$\times \left[ \omega + \frac{\chi_m^0(\zeta^2)}{\tau \chi_m^0(\zeta^2)} \right] \hspace{1cm} (IV, 14)$$

It must be noted that, from eq. (IV, 14), the field threshold does not depend on the wave vector $k$: thus it is independent of sample thickness [5], [17], [18]. This is very different from the conduction regime ($\omega < \omega_c$) where the voltage threshold was independent of thickness. The existence of a field threshold for $\omega > \omega_c$ is well confirmed by the experiments [17].

$\beta$) Spatial periodicity of the molecular pattern at threshold. — At threshold we have

$$E_0^2 = \frac{2}{\lambda} \omega \chi_m^0(\zeta^2, \omega \tau) \hspace{1cm} (i)$$

This is an implicit equation for the wave vector $k$ defining the spatial periodicity. Expanding, as previously, $\chi_m^0(\zeta^2, \omega \tau)$ on powers of $1/\omega \tau$, one finds, for $\omega \tau \gg 1$,

$$E_0^2 = \frac{2}{\lambda} \omega \chi_m^0(\zeta^2) + \frac{1}{\tau} \chi_m^0(\zeta^2) \hspace{1cm} (IV, 15)$$

Inserting eq. (II, 4) and (II, 19), one obtains :

$$\chi_s H^2 + K_{33} \kappa^2 =$$

$$= \frac{2}{\gamma_1} \chi_m^0(\zeta^2) + \frac{1}{\tau} \chi_m^0(\zeta^2) \hspace{1cm} (IV, 16)$$

Eq. (IV, 16) shows that, at threshold,

(i) for fixed $H$, $\kappa^2$ is a linear function of $\omega$

(ii) for fixed $\omega$, the quantity $(\chi_s H^2 + K_{33} \kappa^2)$ remains constant.

These two points have been confirmed by recent experiments on M B B A at Orsay [17].

In order to reach more quantitative comparisons between theory and experiment, we have numerically computed the function $g(X, Y)$ for various values of $\zeta^2 (1 < \zeta^2 < 20)$. On figures (3-4), $1/g(X, Y)$ is plotted versus $1/X$ in the range $X > 5, Y < 2.5$.

A more accurate calculation of $\chi_m^0$ and $\chi_m^0$ has also been carried on a computer for values of $\zeta^2$ corresponding to the experimental conditions.

One finds $\chi_m^0 = 1.04, \chi_m^0 = 0.31$ for $\zeta^2 = 3.05$ and $\chi_m^0 = 0.58, \chi_m^0 = 0.37$ for $\zeta^2 = 4.5$.

A crude estimation of $\zeta^2$ gives 3.2. The slope of the experimental curve $E_0^2(\omega)$ is in a good agreement with the one that can be deduced from the first set $(\chi_m^0, \chi_m^0)$ [17]. The agreement between theory and experiment is less good for the slope of $E_0^2(\omega)$. The slope of this curve depends on the conductivity of the material, and does not remain constant when $\omega$ increases. It appears [17] that the agreement between the theoretical and experimental slopes is reasonable only for small values of $k^2$. The discrepancy that occurs for high values of $k^2$ could be due to the effects of the diffusion currents (see Appendix A).
It must be emphasized that, in this regime, the space charge is time independent, but the curvature (and the molecular preferred orientation) has a rather complicated oscillation with period \(2\pi/\omega_0\). \((\psi(t))^2\) is shown on figure 5. The local dielectric tensor oscillates with the same period. We can then legitimately call this regime a "dielectric regime".

\[
\delta = \frac{2(y_1' - y_2)}{y_1 + 2\beta + 2y_2} = \frac{2\gamma_1}{y_1 - y_2} \frac{1}{1 + \eta_0/\gamma_1}. \quad (V, 3)
\]

**Fig. 5.** \(\psi^2(u)\) is plotted versus \(u\) for : \(u(0, \pi), \zeta^2 = 3.05\) and \(Y = Y_{\infty}^0 = 0.31\).

**V. The case \(\varepsilon_0 > \varepsilon_1\).** — By mixing suitable nematic components one can vary continuously the dielectric anisotropy \(\varepsilon_0\) and even change its sign: in the present section, we shall discuss the instability threshold as a function of \(\varepsilon_0\), allowing \(\varepsilon_0\) to be either positive or negative. We make two restrictions however

\(\alpha\) even when \(\varepsilon_0 > 0\), we consider only the case where the unperturbed optical axis is parallel to the walls (along \(x\), in the notation of this paper), the alignment being due either to a magnetic field \(H\) or to suitable boundary conditions at the walls

\(\beta\) we restrict our attention to nematics with \(\sigma_0 > 0\) (no system with \(\sigma_0 < 0\) being known at present).

In spite of these two restrictions, we shall find a great variety of possible behaviors. The results are summarised in figure 6.

Our investigation will be mainly concerned with the « conduction regime », which will turn out to be often dominant here. The fundamental equation is still eq. (III, 9), which applies for both signs of \(\varepsilon_0\). We may rewrite (III, 9) in a form in which the role of the dielectric anisotropy is more apparent.

\[
\frac{E^2}{D_0} = \frac{1 + \omega^2 x^2}{\varepsilon_0 \omega^2 \tau^2 + \theta_H} \quad (V, 1)
\]

where \(D_0\) is defined by eq. (II, 4), and the Helfrich parameter \(\theta_H\) is given explicitly by:

\[
\theta_H = -\varepsilon_0(\zeta^2 - 1) = \frac{\varepsilon_0}{\sigma_0^0} \left[\sigma_0 \delta + \frac{\sigma_0}{\varepsilon_0} (1 - \delta)\right]. \quad (V, 2)
\]

\(\delta\) is a dimensionless ratio involving the friction coefficients

\[
\delta = \frac{2(y_1' - y_2)}{y_1 + 2\beta + 2y_2} = \frac{2\gamma_1}{y_1 - y_2} \frac{1}{1 + \eta_0/\gamma_1}. \quad (V, 3)
\]

\[
\varepsilon_0^2 = D_0^2 \frac{1 + \omega^2 x^2}{\varepsilon_0 \omega^2 \tau^2 + \theta_H} \quad (V, 1)
\]

From thermodynamic inequalities [16] \(\gamma_1, \eta_0\), and \(\gamma_1 - \gamma_2\) are positive. From the experimental data on alignment by flow in nematics Leslie [13] has concluded that \(-\gamma_2 \geq \gamma_1\). These remarks lead to the inequalities \(0 < \delta < 1\). For \(PAA\) we expect \(\delta \approx 0.7\) (and this order of magnitude also applies probably to \(MBBA\)).

Note that \(\zeta^2\) (in eq. V, 2) may now be positive or negative. Some special values of \(\zeta^2\) play a crucial role in the discussion

\(\alpha\) \(\zeta^2 = 1\). This is obtained when \(\varepsilon_0 = \varepsilon_0^0\), where

\[
\varepsilon_0^0 = -\frac{\delta}{1 - \delta} \frac{\sigma_0}{\sigma_1} e_1 < 0. \quad (V, 4)
\]

For \(\varepsilon_0 < \varepsilon_0^0\) the nematic is absolutely stable under fields of arbitrary strength and frequency (\(\delta\)).

\(\delta\) The value \(\varepsilon_0^0\) defined by (V, 4) is acceptable only if it corresponds to values of \(\varepsilon_0\) and \(e_1\) which are both positive: this requires \(\sigma_0/\sigma_1 < (1 - \delta)/\delta\).
$b) \zeta^2 = 0$. This may occur for two distinct values of $\varepsilon_a$ \(^{(1)}\)

\[
\varepsilon_{a1} = \varepsilon_0 \frac{\sigma_1}{\sigma_0}, \quad \varepsilon_{a2} = \varepsilon_0 \frac{\delta}{\delta} \quad (> 0).
\]

(V, 5)

For $\zeta^2 = 0$ the torques due to the Carr Helfrich effect vanish exactly \(^{(8)}\); eq. (V, 1) gives a threshold independent of frequency $E^2 = D_0^2/\varepsilon_0$. This could have been derived from the standard discussion of the competition between the field $H$ — plus the effect of the walls — (favoring an alignment along $x$) and a static field $\sqrt{E^2}$ (favoring an alignment along $z$).

In the cases which have been studied experimentally we have $\varepsilon_{a1} < \varepsilon_{a2}$. We shall restrict our attention to this case in what follows.

Finally, returning to eq. (V, 1), we find the following types of behavior upon increasing $\varepsilon_a$:

1) $\varepsilon_a \leq \varepsilon_{a0}$ — complete stability

2) $\varepsilon_{a0} < \varepsilon_a < 0$ —
   - conduction regime $\omega < \omega_a$ (this is the case discussed in the previous sections).
   - dielectric regime $\omega > \omega_a$

3) $\varepsilon_a > 0$ conduction regime at all frequencies \(^{(9)}\).

Here we have three sub-classes

a) $0 < \varepsilon_a < \varepsilon_{a1}$ — $E$ increasing function of $\omega$

b) $\varepsilon_{a1} < \varepsilon_a < \varepsilon_{a2}$ — $E$ decreasing function of $\omega$

c) $\varepsilon_{a2} < \varepsilon_a$ — $E$ increasing function of $\omega$

These results are summarised on fig. (6). On this figure we have also delineated the region in which the approximation $\omega T_f \gg 1$ is valid, since this approximation underlies eq. (V, 1). It is clear from the figure that, in case 3), the conduction regime may indeed extend to arbitrarily high frequencies \(^{(9)}\). Some preliminary results on this case have been obtained recently at Freiburg \([19]\).

**VI, Limitations of this model.** — In this model, we have considered the Helfrich-Carr instabilities for an $AC$ electric field in an infinite medium; the boundary conditions at the surface of the sample have been included in a very approximate manner. This limits the validity of the model and may explain some discrepancies with the experimental results. Some other discrepancies may be due to the existence of diffusion currents that we have neglected. Nevertheless it appears that this model is able to explain most of the experimental results for materials having a negative dielectric anisotropy and to give some predictions for the case of a positive dielectric anisotropy.

Provided $d/d_\omega \gg 1$, this model is valid for the whole frequency range of interest ($0 - 10^5$ Hz) except in two small frequency intervals.

a) for very low frequencies ($\omega \leq \omega_0$) where the assumption $\omega T_f \gg 1$ does not hold. However, for very low frequencies, the instabilities may be due to carrier injection (such as in «Felici instabilities») and not to the Carr-Helfrich process;

b) for frequencies of the order of the «cut-off» frequency $\omega_c$. In this range the instability regime is neither a conduction regime nor a dielectric regime, but can be considered as a mixing of both, in agreement with observations made at Orsay.

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**APPENDIX A : Contribution of the diffusion currents.** — We have neglected in section (II, 2) the diffusion currents. We will now take into account these currents and show that their contribution can be neglected.

Let us call $q_i$ the contribution of ion $i$ to the excess-charge and $\overline{D}_i$ its diffusion tensor

\[
\overline{D}_i = D_{ii} + D_{in} (nn) \quad \text{(A, 1)}
\]

The total electric current is

\[
\overline{J} = \overline{J} - \sum_i \overline{D}_i \cdot \text{grad} q_i \quad \text{(A, 2)}
\]

when $\overline{J}$ is defined by eq. (II,7). The charge-balance equation is now

\[
\dot{q} + \frac{q}{\tau} - \sum_i D_{ii} \frac{\partial^2 q_i}{\partial x^2} + \sigma_H E \psi = 0 \quad \text{(A, 3)}
\]

Let us define the average diffusion constant as

\[
D_{\|} = \left( \sum_i D_{ii} q_i \right) / \left( \sum_i q_i \right)
\]

and the Debye screening parameter

\[
q_B^2 = \frac{1}{\tau D_{\|}}
\]

Eq. (A, 3) can be rewritten as

\[
\dot{q} + \frac{q}{\tau} + \sigma_H E \psi = 0 \quad \text{(A, 4)}
\]

with

\[
\frac{1}{\tau} = \frac{1}{\tau} \left( 1 + \frac{k^2}{q_B^2} \right) \quad \text{(A, 5)}
\]

For M B B A, $D \approx 10^{-7}$ cgs; $1/\tau \approx 4 \times 10^2$ s$^{-1}$;
\[ q_D^2 \approx 4 \times 10^9 \text{ cm}^{-2}. \] In the conduction regime, \( k \approx \pi/d \), and taking \( d = 100 \mu \), \( k^2 \approx 10^3 \text{ cm}^{-2} \cdot k^2/q_D^2 \) can be neglected.

In the dielectric regime for frequencies of the order of 400 Hz, one finds \( q^2 \) in the range \( 2 - 4 \times 10^8 \text{ cm}^{-2} \cdot k^2/q_D^2 \) can still be neglected. However it must be emphasized that, for higher frequencies, leading to higher values of \( k^2 \), one could not neglect these terms. In section IV, \( 1/k^2 \) must be replaced by \( 1/\xi^2 (1 + k^2/q_D^2) \). Therefore the results of section IV are no longer valid: \( E^2(\omega) \) and \( E^2(\omega) \) are not straight lines, and, for \( \omega \) fixed, \( \lambda_4 \) is not constant in the range \( k^2 \approx q^2 \).

**APPENDIX B : Extension of the Bloch-Floquet theorem.** — The proof of eq. (II, 32) can be set on the following lines:

(i) The system of equations
\[
\dot{\psi} + f(t) \psi + g(t) z = 0 \\
z + h(t) y + k(t) z = 0
\]

(B, 1)
can be written as
\[
\frac{d}{dt} \hat{\psi}(t) = - (M(t)) \hat{\psi}(t)
\]

(B, 2)
where \( \hat{\psi}(t) \) is a spinor with components \( (y(t), z(t)) \) and \( (M(t)) \) a square matrix
\[
(M(t)) = \begin{bmatrix} f(t) & g(t) \\ h(t) & k(t) \end{bmatrix}.
\]

The general solution is
\[
\tilde{\psi}(t) = a_1 \tilde{\psi}_1(t) + a_2 \tilde{\psi}_2(t)
\]

(B, 3)
where \( \tilde{\psi}_1(t) \) and \( \tilde{\psi}_2(t) \) are linearly independent solutions (and \( a_1 \) and \( a_2 \) are time-independent):
\[
\tilde{\psi}_1(t) = \mathcal{A}(t) \tilde{\psi}_1(0).
\]

Using the Feynman conventions for the ordering of non commuting operators, one can write the evolution operator \( \mathcal{A}(t) \) as
\[
\mathcal{A}(t) = \exp \left( - \int_0^t (M(t)) \, dt \right).
\]

(ii) Let \( (M(t)) \) have period \( T \). If \( \tilde{\psi}(t) \) is a solution of eq. (B, 2), \( \tilde{\psi}(T + T) \) is also a solution of this equation. We can thus write
\[
\begin{align*}
\tilde{\psi}_1(t + T) &= a_{11} \tilde{\psi}_1(t) + a_{12} \tilde{\psi}_2(t) \\
\tilde{\psi}_2(t + T) &= a_{21} \tilde{\psi}_1(t) + a_{22} \tilde{\psi}_2(t)
\end{align*}
\]

where the \( a's \) are time-independent constants.

(iii) We can find two linearly independent spinors \( \tilde{\psi}_1(i = 1,2) \) such as
\[
\tilde{\psi}_1(t + T) = \lambda_i \tilde{\psi}(t) \quad (i = 1,2)
\]

(B, 5)
where \( \lambda_1 \) and \( \lambda_2 \) are the roots of the secular equation
\[
\text{Det} |
\begin{pmatrix}
\lambda & -\lambda d_{ij} & 0 \\
\gamma & \eta & 0 \\
0 & 0 & \gamma + \eta
\end{pmatrix}
\|
= 0
\]

(B, 6)
State
\[
\lambda_i = e^{\nu T} \tilde{\psi}(t)
\]

From eq. (B, 5), the spinor \( \tilde{\psi}(t) \) is periodic with period \( T \). This proves eq. (II, 32).

**APPENDIX C : Torque balance in the conduction regime.** — Let us rewrite eq. (II, 24) as
\[
\frac{\partial \Gamma_{\text{tot}}}{\partial x} = \frac{\partial}{\partial x} (\Gamma_F + \Gamma_H)
\]

(C, 1)
where
\[
\begin{align*}
\frac{\partial \Gamma_F}{\partial x} &= \frac{\gamma_1 \eta_0}{\gamma_1 + \eta_0} \lambda (E^2 - \frac{D_0^2}{\kappa}) \\
\frac{\partial \Gamma_H}{\partial x} &= (\lambda_s H^2 + K_{33} k^2 - \frac{\kappa_0}{4 \pi e_{11}} E^2) \psi
\end{align*}
\]

(C, 2)
and
\[
\frac{\partial \Gamma_H}{\partial x} = \frac{\gamma_1 \eta_0}{\gamma_1 + \eta_0} \frac{\varsigma^2}{\eta}.
\]

(C, 4)
\( \Gamma_F \) is the part of the torque (independent of the space-charge) which is taken into account in the standard discussion of Fredericks transition. \( \Gamma_H \) is the contribution of the Carr-Helfrich process to the torque.

In the «conduction regime», the relaxation time \( T_F \) for the molecular orientation is much longer than the period of the electric field. Hence we may consider only the average values (over one period) of the torque. Then we have
\[
\begin{cases}
\Gamma_F = \frac{\gamma_1 \eta_0}{\gamma_1 + \eta_0} (E^2 - \frac{D_0^2}{\kappa}) \psi \\
\Gamma_H = - \frac{\gamma_1 \eta_0 \lambda}{\gamma_1 + \eta_0} \frac{\varsigma^2}{1 + \omega^2 \varsigma^2} E^2 \psi.
\end{cases}
\]

(C, 5)
From eq. (C, 6), \( \Gamma_H \) vanishes if \( \varsigma^2 \) vanishes.

References

[5] Rondelez (F.), Thesis III cycle, Orsay, 1970; see also ref. [17].