energy interval they can be approximated by ln E and a constant, which yields formula (1), where σ_1 and ν are universal while σ_0 and E_0 depend on the process. The parameter E_0 is determined by the region where the cross section begins to increase. For the K^+p cross section we have $E_0 \simeq 17$ GeV, i.e., several times smaller than for pp. Formula (1) with $\sigma_0 = 17.2$ mb and $E_0 = 17$ GeV, and with the sama parameter values as in (2), describes well the K^+p-cross section growth observed at Serpukhov, and predicts a rather fast growth of this cross section at Batavia (see the figure).

If the total cross section increases like $\ln^2 E$, then it follows from the general axioms [4] that the diffraction-peak slope parameter should increase in accord with the same law, $b \to b_1 \ln^2 E$. It is easy to estimate, however, that the coefficient b_1 is very small. We consider, for example, the rigorous inequalities as $E \to \infty$ (see [4])

$$\sigma \leqslant (4\pi/t_0)\ln^2 E, \quad (d\sigma/dt)_{t=0} \leqslant (\sigma_{el}/4t_0)\ln^2 E, \tag{3}$$

where σ_{el} is the total elastic-scattering cross section. If saturation obtains, then the inequalities (3) can be replaced by equalities, and by virtue of (1) we have $4\pi/t_0 = \sigma_1$. Hence as $E \to \infty$ we have

$$b \equiv (d\sigma/dt)_{t=0}/\sigma_{el} = (\sigma_1/16\pi)\ln^2 E, \tag{4}$$

i.e., $b_1 = \sigma_1/16\pi = 0.04$ $(\text{GeV}/c)^{-2}$. If, as above, we take into account the next terms of the asymptotic expressions, then

$$b = b_0 + 2\alpha'\ln(s/s_0) + b_1[\ln(E/100\text{ GeV})]^2, \tag{5}$$

where the first two terms are the usual ones (in particular, $b_0 = 7.0 \pm 1.2$ $(\text{GeV}/c)^{-2}$) and the last term does not exceed 0.4 $(\text{GeV}/c)^{-2}$ in the 10 - 2000 GeV region, i.e., it is much smaller than the experimental errors.

Thus, the available data do not contradict the universal Froissart growth of the total cross section, and the measurement of the total cross sections in Batavia, particularly the K^+p cross section, is of great interest.

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POSSIBILITY OF LARGE RADIOELECTRIC EFFECT IN A CHOLESTERIC DIELECTRIC

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It is shown that a radioelectric effect, due to the dependence of the pitch of the cholesteric helix on the magnetic field, can be produced in a cholesteric dielectric. The field can be much stronger than in conducting media.

The radioelectric (optoelectric) effect, i.e., the production of a constant electric field E^0 by passage of electromagnetic waves through a medium, has been investigated so far only in conducting media [1 - 3]. We shall show that an analogous effect is possible in dielectric cholesteric liquid crystals, and that if the electromagnetic-wave density has the same value I^0 as in vacuum, then the produced electric field E^0 can be much stronger than in conducting media.

The dielectric tensor ε_{ik} of a cholesteric liquid crystal contains a component that depends on the pitch of the cholesteric helix (see, e.g., [4]); the pitch of the helix depends in turn on the magnetic field H transverse to the helix axis so long as $H < H_{cr}$ (H_{cr} is the field at which the pitch of the helix becomes infinite [5]):

$$\epsilon_{ik} = \epsilon^{I}_{ik} + \epsilon^{II}_{ik}(H^2). \tag{1}$$

In the presence of an external magnetic field $\vec{H}^0$ and an electromagnetic-wave field $\vec{H}'$, the last term of (1) contains a term proportional to $(\vec{H}_0 \cdot \vec{H}')$ (if $\vec{H}_0$ is perpendicular to the helix axis). The presence of this term produces in the electric-induction vector $\vec{D}_i$ a term proportional to $(\vec{H}_0 \cdot \vec{H}')\vec{E}'$ ($\vec{E}'$ is the electric field of the wave), the average of which over the period of the wave differs from zero, and this is indeed the cause of the produced constant radioelectric field. We note that in noncholesteric substances the contribution proportional to $(\vec{H}_0 \cdot \vec{H}')\vec{E}'$ to the induction vector is usually small, owing to the weak dependence of the dielectric constant on the magnetic field. The considered effect is therefore insignificant in noncholesteric substances.

Let the axis of the cholesteric helix coincide with the z axis, let $\vec{H}_0$ be along the x axis, and let I be along the y axis and normal to the sample surface. The flux is so polarized that $\vec{H}'$ is directed along $\vec{H}_0$. We confine ourselves to the wave-frequency region in which the dependence of ε_{ik} on ω can be neglected (from [6] we have $\omega \lesssim 10^5$ sec^{-1}). According to [4], the nonzero components of ε_{ik} are

$$\epsilon_{11} = \epsilon(1 + \delta\cos 2\theta), \quad \epsilon_{22} = \epsilon(1 - \delta\cos 2\theta),$$
$$\epsilon_{33} = \bar{\epsilon} \approx \epsilon, \quad \epsilon_{12} = \epsilon\delta\sin 2\theta \tag{2}$$

where $\delta < 1$ (quantities proportional to δ^2 will be neglected). The angle θ of director rotation is connected with the magnetic field and with the coordinate z by the relations [5]

$$z = \sqrt{K/\chi H^2}\,F(\theta, k), \tag{3}$$

$$k = (2q/\pi)\sqrt{\chi H^2/K}\,E\left(\frac{\pi}{2}, k\right) \tag{4}$$

(q^{-1} is the pitch of the helix in the absence of a field, K is a constant in the expansion of the free-energy density in the derivatives of the director direction, χ is the anisotropic part of the magnetic susceptibility, and $E(\theta, k)$ and $F(\theta, k)$ are elliptic functions of the first and second kind).

Since the electric field of the wave is parallel to the cholesteric axis in the chosen geometry, the angle of rotation of the director is independent of the field. In the wave field we have

$$D_i = \epsilon_{ik}E^0_k + \epsilon'_{ik}E'_k + 2\left(\frac{\partial \epsilon_{ik}}{\partial H^2}\right)(H_0 H')E'_k. \tag{5}$$

Averaging (5) over the period of the wave, expressing H' and E' in terms of I^0, and making use of (2), div $\vec{D}$ = 0, and the absence of a field outside the sample, we obtain (c is the speed of light)

$$E_{\left\{ {x \atop y} \right\}} = \frac{8\pi\delta H_0 I^0}{c\sqrt{\epsilon}(1 + 2\sqrt{\epsilon} + \epsilon)}\,\frac{\partial\theta}{\partial H^2}\begin{Bmatrix}\sin 2\theta\\ \cos 2\theta\end{Bmatrix}. \tag{6}$$

We determine $\partial\theta/\partial H^2$ from (3) and (4). The quantity $E[(\pi/2), k]$ is changed by a factor 1.5 when k changes from 0 to 1. Neglecting this change, we find that $k \approx \text{const}\cdot H$. We then get from (3) and (6)

$$E_{\left\{ {x \atop y} \right\}} = \frac{16\pi\delta I^0\sqrt{1 - k^2\sin^2\theta}}{cH_0\sqrt{\epsilon}(1 + 2\sqrt{\epsilon} + \epsilon)}\,[F(\theta, k) - E(\theta, k)]\begin{Bmatrix}\sin 2\theta\\ \cos 2\theta\end{Bmatrix}.$$

It is seen from (7) that E^0 vanishes as $H_0 \to 0$ and $H_0 \to H_{cr}$, and E^0 is maximal at $H_0 = H_{cr}/2$.

There are two alternatives: 1) measure E^0 with z given, and 2) measure $\langle E^0 \rangle$ (E^0 averaged over the length) in the direction of the helix axis. If the axis spans an integer number of turns, then calculation shows that $\langle E^0 \rangle = 0$. If the number of turns is not integer, then the results of the two alternatives are of the same order of magnitude, equal to

$$E^\circ \approx \frac{32\pi\delta I^\circ}{cH_{cr}\sqrt{\epsilon}(1+2\sqrt{\epsilon}+\epsilon)} \equiv \gamma_\epsilon I^\circ .$$

In conducting media we have, in order of magnitude,

$$E^\circ \approx \frac{\sqrt{2\pi\mu}}{L\omega c}\sqrt{\frac{\sigma}{\omega}}\, I^\circ \equiv \gamma_\sigma I^\circ$$

(μ and σ are the mobility and the conductivity and L is the length of the sample in the flux direction. Here $\sigma \ll \omega$ and both are much lower than the carrier relaxation frequencies. In other cases γ_σ is even smaller.) We see that γ_ϵ is always much larger than γ_σ. We note that the weak damping of the waves in a dielectric makes it possible for E^0 to appear in bulky samples, unlike conducting media in which the field is weak if L exceeds greatly the skin layer.

In conclusion, let us estimate the radioelectric field. At $H_0 \simeq 10^3$ Oe, $\delta \simeq 0.2$, $\epsilon \simeq 4$, and $I^0 = 1$ kW/cm^2 we have $E^0 \simeq 0.1$ V/cm.

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COLLIDING POLES AND GROWING CROSS SECTIONS

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The growth of the total cross sections and other features of pp scattering, observed in high-energy experiments, are discussed on the basis of the assumption of colliding complex Regge poles.

We discuss in the present article the possibility of describing the regularities observed in pp scattering at high energies [1 - 6], namely: (a) the growth of the total cross section $\sigma_{tot}(s)$ in the region $s \sim 500 - 3000$ GeV^2, (b) the change of the sign of $\sigma(s) = \mathrm{Re}T(s, 0)/\mathrm{Im}T(s,$ $\sigma(s) = \mathrm{Re}T(s, 0)/\mathrm{Im}T(s, 0)$ in the region $s \sim 5\times10^2$ GeV^2, and (c) the slowing down of the growth of the diffraction-peak slope parameter b(s) under the assumption [7] that the leading singularities in the j-plane of the partial amplitude of the crossing channel are complex-conjugate poles. Denoting the contribution of the other singularities by F(s, t), we express the scattering amplitude in the form

$$T(s,t) = \beta s^{\alpha} e^{-i\pi\alpha/2} + \beta^* s^{\alpha^*} e^{-i\pi\alpha^*/2} + F(s,t), \quad (1)$$

where $\alpha = \alpha_R = i\alpha_I$ and $\beta = |\beta|e^{i\phi}$ are the trajectory and residue of the complex pole.

We consider first the case when the poles α and α^* collide at t = 0 in such a way that a second-order pole is produced at j = 1.[1)] Since the collision of two poles leads to a root singularity for the trajectory and the residue, we assume that as $t \to 0$ we get $\alpha_I = a\sqrt{-t}$ and $\arg h(t) = b\sqrt{-t}$, where $h = 2i\alpha_I\beta$ and $|h(0)| \equiv h \neq 0$. To describe the low-energy region