that the ratio $r_c/\tau$ is independent of temperature. Therefore, interpretation of the NMR and ESR data on type I stacks does not lead to any contradictions. An estimate of the width of the electron energy band for type I stacks from $\gamma_c$ carried out in the tight-binding approximation gives $0.4$ eV, which is half the corresponding value estimated from the thermoelectric power. This may be an indication of the importance of the enhancement of $\gamma_c$ by the electron-electron interaction. A similar situation occurs also in the case of organic conductors based on TMTSF (Ref. 12).

In the interpretation of the NMR data for type II stacks it is necessary to determine initially whether the electron spins responsible for nuclear relaxation are "pinned" or mobile. In the case of "pinned" spins we have $T_c^* = \tau_c$ (Ref. 8), $\tau_c = \tau_{B}$ is the correlation time of the electron spin. Under the exchange narrowing conditions the ESR line width is $\Delta H = \omega_c/2\gamma_c$, where $\omega_c$ is the second moment of the ESR line system of static spins. In this case the reduction in $\Delta H$ as a result of cooling should be accompanied also by a reduction in $T_c^*$ for type II stacks, which is in conflict with the experimental results. It therefore follows that the model of electron spins "pinned" to type II stacks fails to account for the nuclear relaxation rate data. In other words, the system of nuclear spins should transfer energy to the translational motion of electrons and its dynamics should be described by equations such as Eq. (1). An important piece of evidence in support of this conclusion is that, within the limits of the experimental error, the power exponent $\beta$ governing the low-temperature behavior of $(K^*)_{y}$ is twice as large as $\alpha$. This ratio follows in a natural manner from Eq. (1). Substituting in Eq. (1) the values $\alpha = 0.6$ Oe and $\chi = x_r = x_y$, we find that for parallel orientation of static and mobile spins $\Delta H \approx 60$. This value differs by only 20% from the value of $(K^*)_{y}$ for type I stacks. Consequently, the difference between the rates of relaxation involving stacking of the two types is entirely due to the difference between the contributions of the two types.

The available experimental data are in sufficient to draw the final conclusion on the nature of the interacting electron spins in type II stacks. Bearing in mind an analysis of the various possible models given in Ref. 13 and the nuclear relaxation data, preference should be given to the description of type II stacks within the framework of the model of localization of weakly interacting electrons because of disorder.

A theoretical description suitable for the analysis of the width of the ESR line of type II stacks is not yet available because the dynamics of a one-dimensional spin system with a random interaction is complicated. We can simply note that the reduction of $\Delta H$ we have observed at low temperatures is typical of many disordered organic conductors. It therefore follows that the results of an investigation of the magnetic properties of (TSTf)$_n$H$_2$IO$_4$, by local methods are a demonstration of the magnetic inequivalence of the cationic stacks. The existence of such inequivalent stack makes it possible to explain qualitatively the unusual combination of the transport and magnetic properties of this compound at low temperatures. The problem of the nature of the spins responsible for the rise of the magnetic susceptibility in accordance with a fractional-exponent power law requires further study.

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INTRODUCTION

Many phase transitions are accompanied by an increase in the volume of the elementary cell (see, e.g., Ref. 1). Below such transitions reduce the translational symmetry of the lattice and all the distortions in the low-symmetry phase remain filled, the resulting strong topological interaction between the distortions and the order parameter $\eta$ depends on the symmetry distribution of the latter. The nature of this interaction depends on the symmetry of the order parameter. In symmetric under a continuous group of transformation (in the case, e.g., for an incommensurate charge density wave), a vortex in $\eta$ can be associated with an interaction of such symmetry, a domain wall is associated with the distortion. In the latter case if the crystal contains $N$ parallel distortions, a domain wall structure is in which each wall extends between two distortions. Such a structure has been observed in gadolinium molybdate example.

Impurities and lattice defects can greatly alter the properties of materials near a structural transition. In all cases, effects of distortions on the various quantities characteristic of the phase transition have been found to be proportional to $\eta^2$ rather than $\eta$ owing to the form of a domain wall structure, they are therefore much pronounced at low densities $\eta$. In this paper we consider how distortion domain walls influence the thermodynamic and magnetic properties and their effects on light scattering, the damping of sound and the soft mode for the following type.

EQUILIBRIUM VALUE OF THE ORDER PARAMETER FOR TYPE-II STACKS

We consider a phase transition involving the condensation of a soft mode with a vector $k$ in the Brillouin zone. The displacement field $\Delta r (x) = \eta(x) \exp (ikx)$ is generated in the crystal below the transition temperature, where $\eta(x)$ is a polarization vector for the soft mode and $\eta$ is the order parameter. The Landau expansion for the free energy density valid near the transition point:

$$\frac{f(\eta)}{2} = \frac{a}{2} |\eta|^2 + \frac{b}{2} |\eta|^4 + \frac{c}{4} |\eta|^6$$

Here $a > 0$, $b > 0$, $c > 0$, and $\xi = (\xi_0, \xi_1, \xi_2, \xi_3)$ is the deviation from the mean position of the crystal. When the displacement vector $\eta(x)$ is continuous, it is referred to as the collective coordinate field $\eta(x)$ (Ref. 7).

The soft axis is parallel to the sliding plane and normal to the distortion.

If we neglect deformation effects, which are of little importance in our treatment, the free energy density is again given by (1) in the undeformed coordinates $\eta$. The displacement field generated by the phase transition can be expressed in the form

$$\eta(x) = \eta_0 \exp (i k x)$$

where the function $\eta_0 (x)$ is continuous everywhere except at the cut. On the other hand, in terms of the deformed coordinates $\eta = (\eta_0 \cos k x, \eta_0 \sin k x)$ we can write the displacement field as

$$\eta(x) = \eta_0 \cos (k x + \phi)$$

where $g(x) = \eta_0$ is continuous everywhere. Comparing (3) and (5) we obtain

$$\frac{g(x)}{\eta_0} = \cos \left( k x + \phi_0 \right)$$

Condition (2) and the continuity of $g(x)$ give rise to the boundary condition

$\phi_0 = \pi/2$
The surface of the Washington Monument is capped by a sphere with radius 50 ft. If the height of the monument is 550 ft, what is the slant height of the monument?
where $s = (c_0^2 - \omega^2)/k_0^2$, and $\beta = (\omega/k_0)^2$, and $\omega = (k_0/n)^2$ is the frequency and relaxation time for the $l$th mode. The result (30) shows that the absorption by the domain walls can exceed that due to relaxation when $k_0 > 1$ (i.e., sufficiently far from the transition point).

At low temperatures there is a region in which the inequality $\gamma > k_0 > n/\omega$ holds. Here the absorption is resonant—sound at frequency $\omega$ is absorbed primarily by walls whose dimensions are close to the resonance values ($\omega_m$: $\omega = \omega_m(\omega_m/n)^{1/2}$).

The absorption is greatest at frequencies roughly equal to $\omega_m$. If the $\omega_m$ are soft (so that $\omega_m/n^2$), we shall have $\gamma_0 = 1$ in the region of maximum absorption. Calculations of the absorption using this inequality lead to the result

$$
\eta \approx \frac{1}{\omega_m(n/\omega)^{1/2}} \left(\frac{1}{\gamma_0} \frac{\lambda}{\lambda_0} \right) \frac{\omega_0}{\omega},
$$

where $\omega_0 = 2\omega_m[(k_0/n)^2 - 1]$.

The averaging carried out over the angle $\varphi$ between the projection of the acoustic wave vector and the $z$ axis, over the length $l$ [with the weight function (10)], neglects the anisotropy in $\sigma_{xx}$, $\sigma_{xx}$, and $\sigma_{xx}$. By Fourier transforms, inserting the solution into (25) yields the following expression for the inverse acoustic relaxation time $\tau^{-1}$:

$$
\tau^{-1} = \int \frac{d\omega}{2\pi} \frac{n}{\omega} \int d\varphi \frac{1}{\lambda_0(\omega)} \int d\lambda \frac{\lambda}{\omega_0} \frac{\lambda}{\omega},
$$

where $\omega = (k_0/n)^2$ and $\lambda = \lambda_0/\lambda_0^2$ are the propagation velocity of the wall oscillations and the inverse relaxation time, respectively.

Finally, we obtain $\tau^{-1} = 10^{-11}$. Even at low dissipation densities the walls can thus scatter light much more effectively than the density fluctuations.

5. ABSORPTION OF SOUND

We now consider absorption of a sound wave in a crystal containing dislocation domain walls for which $\lambda > \lambda_0$, where $\lambda$ is the acoustic wavelength. The sound wave causes the walls to vibrate, resulting in the absorption of energy. We examine the simplest case when the interaction energy for the deformation and order parameter is given by

$$
\beta \frac{d^2u}{dx^2} - 2\gamma \frac{du}{dx} + (1 - \gamma) u = 0,
$$

where $\beta = \beta_0 + \beta_1 u$, and $\gamma = \gamma_0 + \gamma_1 u$. The steady-state solution is

$$
\frac{d}{dx} (\beta \frac{d}{dx} u + \gamma_0 u) = 0,
$$

which yields

$$
u = \frac{1}{\beta} \frac{d}{dx} \left( \frac{1}{\gamma_0} \frac{d}{dx} u - u \right),
$$

and

$$
\frac{1}{\beta} \frac{d}{dx} \left( \frac{1}{\gamma_0} \frac{d}{dx} u - u \right) = 0.
$$

The solution is

$$
\frac{1}{\lambda_0} \frac{d}{dx} \left( \frac{1}{\gamma_0} \frac{d}{dx} u - u \right) = 0.
$$

The resulting expression for the inverse acoustic relaxation time is

$$
\tau^{-1} = \int \frac{d\omega}{2\pi} \frac{n}{\omega} \int d\varphi \frac{1}{\lambda_0(\omega)} \int d\lambda \frac{\lambda}{\omega_0} \frac{\lambda}{\omega},
$$

where $\omega = (k_0/n)^2$ and $\lambda = \lambda_0/\lambda_0^2$.

5.1. Dispersion of the Sound Wave

For the propagation of the sound wave in the crystal, we obtain

$$
\frac{1}{\beta} \frac{d}{dx} \left( \frac{1}{\gamma_0} \frac{d}{dx} u - u \right) = 0.
$$

The resulting expression for the inverse acoustic relaxation time is

$$
\tau^{-1} = \int \frac{d\omega}{2\pi} \frac{n}{\omega} \int d\varphi \frac{1}{\lambda_0(\omega)} \int d\lambda \frac{\lambda}{\omega_0} \frac{\lambda}{\omega},
$$

where $\omega = (k_0/n)^2$ and $\lambda = \lambda_0/\lambda_0^2$.

5.2. Dispersion of the Sound Wave

For the propagation of the sound wave in the crystal, we obtain

$$
\frac{1}{\beta} \frac{d}{dx} \left( \frac{1}{\gamma_0} \frac{d}{dx} u - u \right) = 0.
$$

The resulting expression for the inverse acoustic relaxation time is

$$
\tau^{-1} = \int \frac{d\omega}{2\pi} \frac{n}{\omega} \int d\varphi \frac{1}{\lambda_0(\omega)} \int d\lambda \frac{\lambda}{\omega_0} \frac{\lambda}{\omega},
$$

where $\omega = (k_0/n)^2$ and $\lambda = \lambda_0/\lambda_0^2$.
Damping constant for the forced oscillations therefore differs greatly from the damping constant for the eigenmode. In our case the amplitude of the heat source, which is proportional to $\frac{\partial^2 x}{\partial t^2}$, is $\frac{\partial^2 x}{\partial t}$, (34), is spatially constant except within distances $r_0$ from the wall. For $r < r_0$, the temperature distribution is of the form

$$T = \frac{\partial^2 x}{\partial t^2} \exp \left[ -\frac{\partial^2 x}{\partial t^2} \right] \left[ \frac{1}{r_0} \right].$$

Substitution into (35) then yields the estimate

$$\frac{\partial^2 x}{\partial t^2} \approx \frac{2}{\partial^2 x} \left( \frac{\partial^2 x}{\partial t^2} \right)^2 \frac{\partial^2 x}{\partial t^2} \left( \frac{\partial^2 x}{\partial t^2} \right)^2 \left( \frac{1}{r_0} \right)^2 \left[ \frac{1}{r_0} \right].$$

(39)

This result is valid when $\omega \ll \gamma / C_p$, i.e., when the low transfer occurs independently for each wall. At low frequencies $\omega \ll \gamma / C_p$, the departure of the temperature from its equilibrium value is expressible as

$$T = \frac{\partial^2 x}{\partial t^2} \left( \frac{\partial^2 x}{\partial t^2} \right)^2 \left( \frac{1}{r_0} \right)^2 \left[ \frac{1}{r_0} \right],$$

(40)

where the bar denotes a spatial average; $T = \frac{\partial^2 x}{\partial t^2} / C_p$ obeys the equation

$$x \frac{\partial^2 x}{\partial t^2} = -\partial_0 \delta \frac{\partial^2 x}{\partial t^2},$$

where $\delta (\partial_0)$ is the deviation of $\partial_0$ from its average value. We have the estimate

$$\frac{\partial^2 x}{\partial t^2} \approx \frac{2}{\partial^2 x} \left( \frac{\partial^2 x}{\partial t^2} \right)^2 \left( \frac{1}{r_0} \right)^2 \left[ \frac{1}{r_0} \right],$$

where $r_* = r_0 \times \frac{\partial^2 x}{\partial t^2}$. Substituting this into Eq. (35), we obtain

$$\frac{\partial^2 x}{\partial t^2} \approx \frac{2}{\partial^2 x} \left( \frac{\partial^2 x}{\partial t^2} \right)^2 \left( \frac{1}{r_0} \right)^2 \left[ \frac{1}{r_0} \right].$$

(41)

At low frequencies the effective damping constant for the forced oscillations is thus independent of the dislocation density. For $T = T_0$ and $n_0 \approx 10^{10}$ cm$^{-2}$, the transition from (39) to (41) occurs at frequencies $\omega = 10^{-10}$ s$^{-1}$.

7. CONCLUSIONS

We have used a simple model to analyze how dislocations affect the properties of phase transitions that are accompanied by an increase in the volume of the elementary cell. The model treats a phase transition of the dislocation type in which a single soft mode condenses. The situation for real crystals is frequently more complicated. For example, in lead zirconate (an antiferroelectric), several soft modes condense simultaneously, and the elementary cell for the low-temperature phase contains eight elementary cells for the high-temperature phase. In nonferroelectric ferroelectrics, the order parameter has two components. There are also materials (such as ammonium dihydrogen phosphate, ADP) in which the period increases as a result of an order-disorder transition. Although the specific models for the interaction of the order parameter with dislocations differ, in all cases dislocations are responsible for domain wall formation when $T = T_0$, and the results found above remain qualitatively correct. Above the transition point, the order parameter interacts with the dislocations either through the deformation field or because the order parameter is distorted at the centers of defects. For $T > T_0$, all the effects involved in the interaction are proportional to the dislocation density $n_0$, while for $T < T_0$ they are proportional to $n_0^2$, owing to the formation of domain walls in this case. It would be extremely interesting to have comprehensive experimental data on how plastic deformation alters the properties of phase transitions that increase the volume of the elementary cell.

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Translated by A. Mason