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Matter in a Superstrong Magnetic Field

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The physical properties of atoms, molecules, and solids in a superstrong magnetic field up to 10^{12} to 10^{13} Oe are considered in the quantum "drift" approximation. It is shown that the atoms, which are strongly elongated along the magnetic field, merge together into polymer molecules at low temperatures and, therefore, a solid should have the structure of a polymer.

1. INTRODUCTION

ACCORDING to contemporary ideas,[1] magnetic fields up to 10^{12} to 10^{14} Oe may exist in neutron stars or on their surfaces. Such fields must strongly influence the physical properties of matter. In fact, the interaction of atomic electrons with an external magnetic field becomes larger than their Coulomb interaction if the magnetic field exceeds the magnitude $m^2e^3c\hbar^{-3} = 2.35 \times 10^9$ Oe (in what follows the magnetic field B will be measured in precisely these units, and all remaining quantities will be expressed in atomic units). Therefore a complete rearrangement of an atom's electron shells must occur for $B \gg 1$. In particular, for the hydrogen atom with $B \gg 1$ the electron cloud takes the shape of a thin needle, which is elongated along the direction of the magnetic field.[2–4] For heavy atoms with atomic number $Z \gg 1$ the spherical symmetry is preserved for $B < Z^3$,[5] but then for $B > Z^3$ they are also elongated along the magnetic field.[6,7] Markedly elongated atoms possess a large electric quadrupole moment, and at not very high temperatures, the forces of attraction acting between these moments should lead to the formation of molecules[8] and their condensation into a solid phase.

In the present article we consider certain physical properties of matter in a superstrong magnetic field, starting with the interaction of isolated atoms and ending with the fundamental characteristics of the solid state.

2. THE QUANTUM "DRIFT" APPROXIMATION

In a very strong magnetic field, that is (in units of $m^2e^3c\hbar^{-3}$) for $B \gg 1$, in the excited state all of the electrons must be in the very lowest Landau level, and the spin magnetic moments are oriented along the field. In this connection each electron executes oscillations along **B**, but in the plane perpendicular to **B** it moves in a narrow cylindrical shell whose width is equal to the Larmor radius $\sim 1/\sqrt{B}$. In classical mechanics the so-called drift approximation is used in order to describe such motion in strong magnetic and slowly varying electric fields. It is natural to utilize the analogous approximation in the quantum case.

The Schrödinger equation for electrons with spin $\sigma = -1/2$ (in atomic units) has the form

$$\tfrac{1}{2}(\mathbf{p}+\mathbf{A})^2\psi - \varphi\psi - \tfrac{1}{2}B\psi = E\psi, \tag{1}$$

where $\hat{\mathbf{p}} = -i\nabla$ and φ is the potential of the electric field. Let us choose the vector potential **A** in the form

$$\mathbf{A} = b\nabla\chi, \quad \mathbf{B} = [\nabla b \nabla\chi]; \tag{2}^*$$

the z axis is directed along **B**.

We shall assume that the functions b(x, y) and $\chi(x, y)$ do not depend on z. The arbitrariness in their choice is restricted only by the second relation in (2). Employing this relation, we select the quantity b in such a way that the lines b = const correspond to the lines of drift of the electrons in the plane perpendicular to **B**. In addition, we normalize χ so that it varies from 0 to 2π (this can be achieved by the introduction of a new function $b_1 = b_1(b)$ in place of b, which is constant on the same lines b = const).

Let $\psi = \psi_m = e^{-im\chi}\psi'$, where $\psi'(b, \chi)$ is a periodic function of χ, and m is the azimuthal quantum number. It is not difficult to verify that ψ' satisfies the same equation of the form (1), with **A** replaced by $\mathbf{A}' = \mathbf{A} - m\nabla\chi = (b - m)\nabla\chi$. We shall seek the quasi-classical solution with $m \gg 1$, which is localized near b = m and which varies slowly with respect to χ along the line b = m, so that $|\partial\psi'/\partial\chi| \ll m|\psi'|$. In accordance with this, one can neglect the derivatives of ψ' with respect to χ in Eq. (1). In addition, in the zero-order approximation of the expansion in powers of 1/B, one can neglect the small terms $(1/2)p_z^2\psi' + \varphi\psi'$, so that Eq. (1) for ψ' takes the form

$$-\frac{1}{2}\frac{\partial^2\psi'}{\partial\xi^2}(\nabla b)^2 + i(\nabla b\nabla\chi)\left(\xi\frac{\partial\psi'}{\partial\xi} + \frac{1}{2}\psi'\right) + \frac{1}{2}(\nabla\chi)^2\xi^2\psi' - \frac{B}{2}\psi' = E_0\psi', \tag{3}$$

where $\xi = b - m$, $B = [(\nabla b)^2(\nabla\chi)^2 - (\nabla b\cdot\nabla\chi)^2]^{1/2}$ and in addition we omitted the small term $-(1/2)(\partial\psi'/\partial\xi)\Delta b$. Equation (3) has the solution

$$\psi' = \exp\left(-\frac{C\xi^2}{2}\right), \quad C = \frac{B - i(\nabla b\nabla\chi)}{(\nabla b)^2} \tag{4}$$

with the eigenvalue $E_0 = 0$. As we see, this solution corresponds to a wave function localized near the line b = m = const with the width of the localization in ξ of the order of $1/\sqrt{B}$. One can interpret it as the function describing the electron's drift along the line b = m in the plane perpendicular to **B**.

We shall assume that the lines of drift b(x, y) = const

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*$[\nabla b\nabla\chi] \equiv \nabla b \times \nabla\chi$.

are closed; then these closed lines of drift form a certain family of lines enclosed one inside the other. Near the center (b = 0) let the lines b = const be close to circles, so that for small values of b the functions $\sqrt{2b/B}$ and χ coincide with the cylindrical coordinates ρ and θ. Then for small values of m the wave functions will be close to functions having a constant angular momentum, $\psi_m = \rho^m \exp(-B\rho^2/2)$. For sufficiently large values of m these functions approach $\exp[-B(\rho - m)^2/2]$ and, as we see, go over into the quasiclassical functions (4) with $(\nabla b \cdot \nabla\chi) = 0$.

Thus, for a sufficiently broad class of functions b, coinciding with $B\rho^2/2$ for small b and smoothly varying at large distances, one can choose the wave functions in the form of "drift shells." In the zero-order approximation, all of these functions have one and the same eigenvalue, $E_0 = 0$. The found functions are approximately orthogonal to each other: for a large difference $m - m'$ they simply do not overlap with respect to b, and for small values of $m - m'$ their dependences on b are similar and they are orthogonal due to their angular parts $e^{-im\chi}$.

Since the level E_0 is infinitely degenerate, then in order to obtain the next approximation it is necessary to solve the secular equation. Let us choose the function b such that only the diagonal terms remain in the secular equation; then it takes the form

$$-\frac{1}{2}\frac{\partial^2\psi_{ml}}{\partial z^2} - \varphi_m\psi_{ml} = E_{ml}\psi_{ml}, \tag{5}$$

where φ_m represents the potential of the electric field, averaged with the m-th wave function of the variables b and χ, and here the function ψ_{ml} only depends on z; l corresponds to the longitudinal quantum number.

Now it is not difficult to see how one should choose the function b. Since according to Eq. (5) the Hamiltonian of the longitudinal motion is diagonal in the functions ψ_{ml}, this means that for quasiclassical functions with $m \gg 1$ the value E_{ml} along the drift shell b = const remains constant. In other words, for each l the functions b are chosen from the condition of constancy of the energy eigenvalue E_{ml} of the longitudinal motion as a function of x and y. This condition is analogous to the condition for the conservation of the longitudinal adiabatic invariant that determines the drift shells in the classical case. Even in the case of small values of m, the condition that the Hamiltonian be diagonalized has a similar meaning.

The arguments given here primarily pertain to the assumed case when axial symmetry does not exist. However, in the presence of such symmetry the entire picture is much simpler, since here $b = \rho^2/2$ and the radial dependence of ψ_m is exactly known.

We note that the relation $B = \nabla b \times \nabla\chi$ implies that the Jacobian of the transformation from the variables b, χ to the variables x, y is equal to B. From here it follows that $2\pi b = BS$, where S denotes the area bounded by the curve b = const. Thus one can say that the area associated with the drifting motion is quantized: In connection with a transition from one azimuthal level m to another, the area increases by the amount $S_0 = 2\pi/B$. One would also anticipate this result according to the correspondence principle, since in the classical case the area S for B = const corresponds to the third adiabatic invariant. Therefore, during the interaction of elongated atoms ($B \gg Z^3$) with external fields and with other atoms, the area of their transverse cross section remains constant and only its shape changes.

3. ATOMS IN A SUPERSTRONG MAGNETIC FIELD

Many-electron atoms in the presence of a superstrong magnetic field have been considered in[5-7]. It was shown that for $B < Z^3$, in the ground state several electrons are found in each m-level, and a modified Thomas-Fermi model may be used to describe the ground state. In this approximation the atom keeps its spherical symmetry, but its characteristic radius varies like $Z^{1/5}B^{-2/5}$. For $B > Z^3$ only one electron remains in each m-level, and for $B \gg Z^3$ the atom is strongly elongated along the direction of the magnetic field, taking the shape of a needle of radius $\rho_z = \sqrt{2Z/B}$. In this connection all of the electrons are found in the ground level of the longitudinal motion, whose depth increases without limit with increasing magnetic field, and for $B \gg Z^3$ becomes considerably larger than the Coulomb interaction energy at the average distance from the nucleus. Here the wave function can be approximated by the exponential $\psi = \exp(-\alpha|z|)$, and the binding energy of the atom, after minimization with respect to α and neglecting the exchange corrections, turns out to be given by

$$E_a = -\frac{9}{8}Z^3L^2, \quad L = \ln\frac{1}{\alpha\rho_z} \cong \frac{1}{2}\ln\frac{B}{Z^3}. \tag{6}$$

This expression can be written in the form $E_a = Z\epsilon$, where ϵ is the average binding energy of a single electron, which in fact corresponds to the binding energy of an electron on the periphery, that is, for $\rho = \rho_z = \sqrt{2Z/B}$, since upon averaging it is precisely these electrons which give the major contribution.

Expression (6) determines the energy of an atom at rest. For an atom moving with velocity **V**, we must add to expression (6) the kinetic energy $MV^2/2$ (M denotes the mass of the atom, which is essentially the same as the mass of the nucleus) and the potential energy of the deformation of the electron shell of the moving atom. One can find this energy by taking into account the properties of the drifting motion of the electrons which was considered above.

In fact, if we change to a coordinate system moving together with the atom along the x axis, then in this system there is a transverse electric field $E_y = -V_\perp B$, whose presence leads to such a displacement of the drift surfaces so that the energy on a drift line would be constant. Thus, for small $V_\perp$ the displacement ξ along the radius is obtained from the condition

$$\frac{\partial\varepsilon}{\partial\rho}\xi - V_\perp B\rho\sin\vartheta = 0. \tag{7}$$

Since as usual the transverse area $S_0 = 2\pi/B$ occurs for each electron, then the displacement across **B** is incompressible, and in this connection the energy of all the internal electrons does not change—the entire change of energy is due to the surface electrons, which are shifted farther from the nucleus. This change is given by

$$\delta E = \frac{1}{S_0}\int\left(\frac{\partial\varepsilon}{\partial\rho}\frac{\xi^2}{2}\right)_{\rho=\rho_z}\rho\, d\vartheta, \tag{8}$$

where $1/S_0$ denotes the density of electrons. Substituting here ξ from (7) and taking into consideration that

$$\varepsilon=-\frac{9}{8}Z^2L^2,\quad \frac{\partial\varepsilon}{\partial\rho}=\varepsilon\frac{2}{L\rho},\quad \rho_z{}^2=\frac{2Z}{B},$$

we obtain $\delta E = (4/9)BV_{\perp}^2/L$ so that the total energy is given by

$$E=E_a+\frac{MV^2}{2}+\frac{M_0V_{\perp}{}^2}{2},\quad M_0=\frac{8}{9}\frac{B}{L}. \tag{9}$$

For large B the quantity M_0 may be comparable with M, and therefore an appreciable anisotropy of the atom's effective mass appears: $M_{\parallel}^* = M$, $M_{\perp}^* = M + M_0$. The quantity M_0 was determined in[9] for a hydrogen-like atom (for a Mott exciton). In our notation it is given by $M_0 = B/2L$. This expression differs from (9) only by a numerical factor, and can be obtained from (8) by substituting the value $\epsilon = -2L^2$ for the hydrogen atom into it.

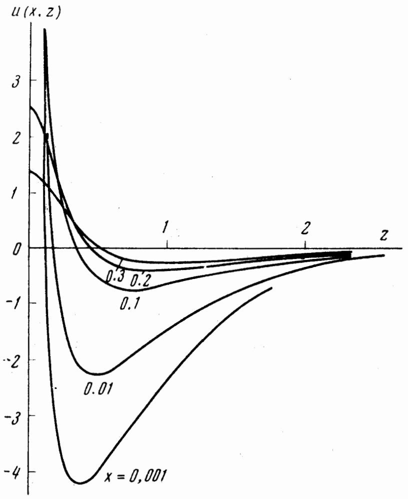

FIG. 1.

4. THE INTERACTION OF ATOMS; MOLECULES

The strongly elongated atoms associated with $B \gg Z^3$ possess large electric quadrupole moments and therefore must interact strongly among themselves. For a large distance r, two atoms interact simply as two quadrupoles, which are rigidly oriented along the direction of the magnetic field. For the distribution $n \propto \exp(-2\alpha|x|)$ of the electron density, the atom's quadrupole moment is given by $Q_{zz} = 2\langle z^2\rangle = \alpha^{-2}$ and for $\alpha r \gg 1$ the interaction energy U is given by

$$U=\frac{9Z^2}{2\alpha^4r^5}P_4(\cos\theta), \tag{10}$$

where θ is the angle between the direction of the magnetic field and the vector joining the centers of the two atoms, and $P_4(x) = (1/8)(35x^4 - 30x^2 + 3)$ is the Legendre polynomial.

The potential (10) is a monotonic function of r; therefore the sign of the interaction force is determined by the sign of the polynomial P_4: when $P_4 > 0$ the atoms repel each other, and when $P_4 < 0$ they attract. Consequently repulsion occurs for $\theta = 0$ and $\theta = \pi/2$, and attraction occurs in a certain range of angles near $\theta = 49°$, where $\cos^2\theta = 3/7$ and P_4 has a minimum.

For $\alpha r \leq 1$ the interaction energy of needle-shaped atoms with rigidly fixed electron density per unit length $\propto \exp(-2\alpha|z|)$ can be found numerically. One can represent the potential U in the form $U = Z^2\alpha u(x, z)$, where x and z are the dimensionless components of the vector joining the centers of the atoms, that is, $x = \alpha a$ and $z = \alpha c$, where a denotes the distance between the axes of the atoms and c denotes the distance between the nuclei along the direction of the field **B**. The potential u(x, z) as a function of z is depicted in Fig. 1 for several values of x. It is seen that u(x, z), for a fixed value of x, has a minimum $u_m(x)$ whose value as a function of x is depicted in Fig. 2. As is seen, $u_m \sim 1$, and since $\alpha \sim 1/ZL$, then the interaction energy U is at most L times smaller than the total binding energy of each of the atoms.

Since $u_m(x)$ is a monotonic function of x, the interacting atoms will approach each other until they come into contact. Upon further transverse motion of the nuclei toward one and the same line of force, the electrons will become collectivized into external shells whose drift trajectories will encompass both nuclei. One can say that these electrons are "extruded" from the more interior shells, since the electron shell is incompressible in a direction transverse to the magnetic field. This effect, which leads to a certain average removal of the electrons from the nucleus and to an increase of the potential energy, is equivalent to the appearance of repulsive forces upon contact of the atoms. If L is not large, so that contact begins for rather large values of x, when the value of U is small according to Figs. 1 and 2, then a molecule having a small number of collectivized electrons will be formed out of the two atoms which are stuck together. On the other hand, for large values of L a kind of supermolecule is formed in which all of the electrons are collectivized and form a single needle-shaped cloud along whose axis the nuclei are distributed.[8] In a very strong field such supermolecules can be formed out of several, let us say n, atoms; in this connection the nuclei are distributed at a certain distance R from each other along a single line of force. First let us consider precisely such molecules.

The radius of the electron cloud in such a molecule is obviously given by $\rho_{nZ} = \sqrt{2nZ/B}$. In a very strong magnetic field, when the molecule becomes extremely narrow and the motion of the electrons is essentially one-dimensional, the wave function of the longitudinal motion will correspond to a bound state in a potential having the form of a superposition of one-dimensional δ-functions at the points where the nuclei are located.

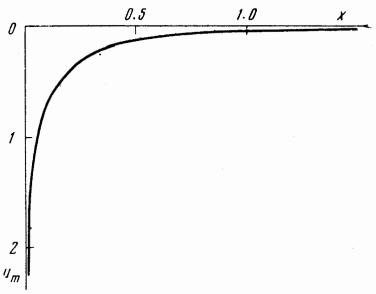

FIG. 2

On the outside away from the nuclei this function has the form $\exp(-\alpha|z|)$, and between the nuclei it has the form $\cosh \alpha z$. But if the nuclei are situated sufficiently close to each other, so that $\alpha R \ll 1$, then one can neglect the "dip" in the function ψ between nuclei and set $\psi = \text{const}$. If we approximately assume α to be identical for all electrons, then the corresponding electronic density per unit length, λ, has the form $\lambda = nZ\beta\exp(-\alpha|z|)$ outside of the nuclei and $\lambda = nZ\beta = \text{const}$ between the nuclei. From the normalization condition it follows that $\beta R(n-1) + \beta/\alpha = 1$. In this approximation and with logarithmic accuracy, the energy nE_n of an n-atom molecule (E_n denotes the energy associated with a single atom) is determined by the following expression

$$nE_n = nZ\left\{\frac{\alpha\beta}{2} - 2nZL_n\beta - \frac{1}{2}nZL_n\beta + \frac{1}{2}(n-1)nZ\beta^2R\right\} + nf_n\frac{Z^2}{R}, \tag{11}$$

where

$$L_n = \ln(1/\beta\rho_{nz}), \quad f_n = {}^1/_2 + {}^1/_3 + \cdots + 1/n.$$

The first term inside the curly brackets in Eq. (11) represents the kinetic energy of the electrons, the second term represents the interaction energy of the electrons with the nuclei, and the last two terms represent the Coulomb interaction energy of the electrons among themselves. The last term in (11) corresponds to the interaction energy of the nuclei. One should regard the quantities β and R in (11) as free parameters, which must be determined by minimizing the energy E_n. Since the minimum with respect to β is reached for $\beta \sim nZ$, then to within logarithmic accuracy we have $L_n \cong (1/2)\ln[B/(nZ)^3]$. To within the same degree of accuracy, in minimizing expression (11) with respect to β and R one can regard the quantity L_n as constant.

In terms of the variables $\xi = \beta R(n-1)$ and $\eta = \beta/nZL_n$ the expression for E_n reduces to the simpler form

$$E_n = -\frac{1}{2}n^2Z^2\left\{\frac{\eta^2}{1-\xi} - 3\eta + \xi\eta + A_n\frac{\xi}{\eta}\right\}, \tag{12}$$

where $A_n = 2f_n(n-1)/nL_n$. From here it follows that the minimal value of E_n with respect to ξ and η has the form $E_n = -n^2Z^2L_n^2F(A_n)$, where one can easily determine the function $F(A_n)$ which depends on a single variable by minimizing expression (12). This function decreases with increasing A_n, that is, the molecule's energy E_n increases in absolute magnitude faster than L_n^2 with increasing magnetic field, and with increasing B the formation of larger and larger molecules becomes favorable. However, since the dependence on B is logarithmic, the corresponding values of the magnetic field turn out to be very large. Namely, the calculation shows that $E_2 < E_1$ only for $L \equiv (1/2)\ln(B/Z^3) > 5$ and $E_3 < E_2$ for $L > 6$, that is, the formation of diatomic molecules, and all the more for triatomic and many-atom molecules with completely collectivized electrons, only occurs for very large values of B.

From here it follows that for $B \gg Z^3$, but not too large values of the magnetic field, the formation of the molecules must occur by means of the "adhesion" of the atoms to each other with the preservation of their "individuality." The adhesion energy of two atoms in a molecule is determined by the quantity u_m. If a third atom is brought up to such a diatomic molecule, then it becomes attached to it with roughly the same binding energy, since u(x, z) rather rapidly decreases with increasing z. In this connection, since the atoms repel each other for small values of z (see Fig. 1), then the third atom attaches itself to the external side of the molecule (along the magnetic field). The subsequent atoms will attach themselves to the ends of the molecules in analogous fashion, so that a polymer thread will be formed when the number of atoms is large. The binding energy of an individual atom in such a thread is of the order of magnitude u_m, that is, it is in all only L times smaller than the total binding energy of an isolated atom.

5. THE SOLID STATE

At sufficiently low temperatures the attractive forces between the atoms due to the presence of their large electric quadrupole moments must lead to their condensation into the solid state. At first glance it appears that in connection with the condensation of the "needle-shaped" atoms which are markedly elongated along the magnetic field, nothing hinders them from being packed alongside of each other. Here the electrons must occupy all space with more or less uniform density, and by virtue of their Coulomb repulsion the nuclei submerged in the electronic liquid must form a crystal lattice structure of the close-packing type. It is obvious that in such condensation, the average distance between the nuclei turns out to be much smaller than the length of the isolated atoms out of which the solid was formed. In this connection the characteristic de Broglie wavelength of the longitudinal motion of the electrons must be much larger than the average distance between nuclei. In other words, with respect to the longitudinal motion we have, as it were, a strongly compressed substance. In this connection, by virtue of the smallness of the wave vector k we can approximately replace the periodic function $v_k(z)$ by the function $v_0(z)$ corresponding to $k = 0$ in the Bloch functions $\psi_k = v_k(z)e^{ikz}$ of the longitudinal motion. Here the part of the kinetic energy of the electrons, corresponding to the exponential factor e^{ikz}, is determined as for the case of a free electron gas and is given by $\epsilon = \langle p_z^2/2\rangle = (2/3)\pi^2N_{\max}^2$, where $N_{\max}$ denotes the limiting value of the occupation number of the levels of longitudinal motion per unit length.

Now let us take into consideration that in a superstrong magnetic field, in the ground state all of the electrons are found in the lowest Landau level, and to each state there corresponds a small area in the plane perpendicular to the magnetic field, this area being equal to $S_0 = 2\pi/B$. Since only a single electron is found in each of the possible states, then $N_{\max} = Zn/S_0 = 2\pi Zn/B$, where n denotes the density of nuclei, and nZ denotes the density of electrons. In this approximation and under the assumption that only one nucleus occurs per elementary cell, the total energy per atom is determined by the expression

$$E = \frac{8\pi^4Z^3n^2}{3B^2} + Z\int_\Omega \frac{1}{2}\left(\frac{\partial v_0}{\partial z}\right)^2 d\mathbf{r} + \frac{Z}{2}\varphi_i - \frac{Z-1}{2}\int_\Omega \varphi v_0{}^2 d\mathbf{r}, \tag{13}$$

where the integration is carried out over the volume Ω

of the elementary cell, φ is the potential of the electric field, φ_i is the potential at the point nucleus created by all of the charges except that of the given nucleus, and v_0^2 is normalized to unity. The second term in expression (13) corresponds to the kinetic energy of the longitudinal motion. We have neglected the exchange interaction, considering that $Z \gg 1$. To this same degree of accuracy, one can replace $Z - 1$ by Z in the last term.

The first term in (13) is proportional to n^2, and for fixed configurations of the charge distribution the second and the last two terms are proportional to $n^{2/3}$ and $n^{1/3}$, respectively, that is, they depend on n considerably more weakly. Therefore, one can approximately seek the minimum of E with respect to v_0, so that for the minimum value of E with respect to n we have

$$E = 8\pi^4 Z^3 n^2 / 3B^2 - AZ^2 n^{1/3}. \tag{14}$$

Here A is a certain constant of the order of unity, which very weakly depends on n, so that one can seek the minimum of expression (14) with respect to n for constant A. From the condition $\partial E / \partial n = 0$ we obtain the following expression for the minimum value of E:

$$E = -\frac{5}{6}\left(\frac{B\sqrt{A}}{4\pi^2 Z^3}\right)^{2/5} AZ^3. \tag{15}$$

One can approximately estimate the quantity A by replacing the elementary cell by a sphere of radius $R = (3/4\pi n)^{1/3}$ and by choosing the function v_0 to be spherically symmetric, $v_0 = v_0(r)$. For the simplest trial function $v_0 = \text{const}$, corresponding to the free electron gas, the quantity A is equal to $(9/10)(4\pi/3)^{1/3} = 1.45$. In actual fact, however, the electronic density must be several times larger around the nucleus, that is, v_0 must be a decreasing function of r. However, even a trial function of the form $v_0 = \text{const} \cdot (1 - r/R)$ which decreases strongly with r, and corresponds to a vanishing of the electron density on the boundary of the cell, increases the value of the constant A by only a few percent for the value $B = Z^3$ and doesn't change its value for $B = 2Z^3$. For $B > 2Z^3$ the free-electron approximation, $v_0 = \text{const}$, must be even better satisfied, so that for all $B > Z^3$ the binding energy of the solid state, under the assumption of close packing, is given sufficiently exactly by expression (15) with $A = 1.45$.

But from here one can immediately reach the conclusion that in actual fact the density of packing in uncompressed matter will not be attained. In fact, as one can easily verify, the quantity (15) for $A = 1.45$ lies above the binding energy of a diatomic molecule with completely collectivized electrons $E_2 \cong (9/2)Z^3L_2^2$, where $L_2 = (1/2)\ln[B/(2Z)^3]$, and (15) is comparable with E_2 only for unrealistically large values of the magnetic field, $B/Z^3 \sim 10^7$. Thus, close packing is energetically unfavorable even in comparison with a gas of diatomic molecules, and it is all the more unfavorable in comparison with polymer filaments. The formation of a solid by means of the adhesion of polymer molecules is more natural. It is rather difficult to estimate the binding energy of such adhesion, but apparently it is appreciably smaller than the binding energy of a single atom in the molecule.

Thus, upon heating the substance one would first expect the breaking of the bonds between the polymer molecules, then side by side with the breaking of the molecules into fragments and the reduction of their average lengths will also occur the excitation of electronic m-levels having energies of the order of ZL, and only in the final stage will ionization and dissociation of the molecules into atoms occur with an energy dissipation $\sim Z^3L$.

CONCLUSION

Thus, the semi-quantitative investigation carried out above indicates that a substantial change in the physical properties of matter occurs in the presence of a superstrong magnetic field. First of all, the dimensions of the atoms are reduced and the binding energies and ionization potentials are increased. For $B \gg Z^3$ the atoms are markedly elongated along the direction of the magnetic field, and the atoms begin to strongly interact between themselves, sticking together into long polymer molecules. At sufficiently low temperatures these molecules can condense into the solid state, so that the structure of the solid must have the character of a polymer; estimates show that the usual crystalline lattice with closepacking turns out to be energetically unfavorable. In addition to the ground state which we have treated here, the investigation of excited atoms and molecules is of interest.

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