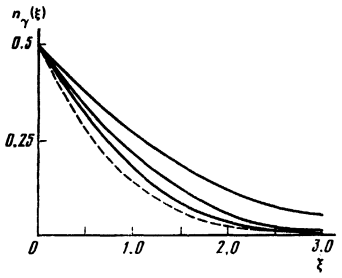

FIG. 2. Plot of the distribution $n_\gamma(\xi)$ at $\gamma=1$, 1.5, and ∞ (the solid lines, reading down); dashed—temperature approximation for $\gamma=\infty$.

consideration that $f(\xi/\tilde{\Theta})$ is that limit to which $n_\gamma(\xi)$ must tend in the case of a "gradual" turning on of the interactions between the electrons and $\mathscr{P}$ is constant, so that in final analysis the thermalization condition $\nu_{ee} \gg \nu_{ep}^{\varepsilon}$ is satisfied. Calculation of the field contribution to the resistivity in the temperature approximation leads to the equation

$$\delta\tilde{\rho}_{ph}{}^{T}(E)=\frac{4(\gamma^5-1)}{5D_5}\rho_{ph}(T). \tag{10}$$

Comparison of (3) and (10) as functions of the parameter γ shows that the difference between them in the entire investigated interval $1.1<\gamma<10$ is less than 1%.

Thus, for the considered simplest model of electrons with isotropic and quadratic dispersion law, and phonons with a Debye spectrum, in the absence of collisions between the electrons, the temperature approximation describes an integral characteristic [that depends on the distribution function $n_\gamma(\xi)$] such as $\delta\rho_{ph}^{T}(E)$ with very high accuracy. The possible cause of this accuracy, in our opinion, may be that both $\mathscr{P}$ and $\delta\rho_{ph}^{T}(E)$ are, in the model investigated by us, actually moments of the same order of the function $n_\gamma(\xi)$. Therefore, although the functions $n_\gamma(\xi)$ and $f[\xi/\tilde{\Theta}(\gamma)]$ differ insignificantly (but still noticeably, see Fig. 2), the proximity of $\delta\tilde{\rho}_{ph}(E)$ to $\delta\rho_{ph}(E)$ can be of quite high order, since $\tilde{\mathscr{P}}_\gamma = \mathscr{P}_\gamma$ by definition. This in turn leads to the assumption that if we consider in the temperature approximation defined above the physical responses corresponding to moments of $n_\gamma(\xi)$ of a different order than $\mathscr{P}_\gamma$, then the difference between these responses and the "true" ones should be more noticeable.

We remark also that by combining Eqs. (6) and (10) we see that from the connection between the experimentally observed quantities[3] $\delta\rho_{ph}^{T}(E)$ and $\mathscr{P}$, the electron-phonon interaction constant[1] α drops out in the temperature approximation. Namely, a linear connection exists of the type

$$\delta\tilde{\rho}_{ph}{}^{T}(E)=4\mathscr{P}/3j_0{}^2. \tag{11}$$

In conclusion, the authors thank A. A. Motornaya for great help with the computer calculations.

[1)]We note, however that we were unable to obtain the exact form of $a(\gamma)$, primarily because of the nonlocality and nonlinearity of the investigated equation.

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Infrared absorption of light by small-scale fluctuations in compensated semiconductors

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We calculate the absorption coefficient of light in transitions between energy levels of small-scale fluctuations of impurities in weakly doped compensated semiconductors. The absorption spectrum is hydrogenlike. The shape of the absorption line corresponding to a transition from the ground state to the first excited state of the fluctuations is calculated. It is shown that the line width is due to deviation of the fluctuation shape from spherical.

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The far tail of the state density in doped semiconductors is determined by the small-scale fluctuations whose potential coincides with the potential of the effective point charge $Z>1$.[1] Such a system has hydrogenlike energy levels. This paper discusses the possibility of direct optical observation of transitions between the energy levels of a small-scale fluctuation. The frequency range of the investigated transitions lies in the far infrared or in the submillimeter region.

To observe the transition between the levels of such a small-scale fluctuations it is necessary that the state

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density near the Fermi level be determined precisely by these small-scale fluctuations. This situation is realized in weakly doped and strongly compensated semiconductors with a correlated arrangement of the impurities. The correlation is due to repulsion between the impurities at the temperature at which the sample is prepared.[1] In this case the radius of the large-scale fluctuations is limited by the ion screening radius $r_i = (\varkappa T_0/8\pi Ne^2)^{1/2}$ where T_0 is the diffusion-quenching temperature.

As shown in Ref. 1 (p.342), at energies higher than E_{c_1}, where

$$E_{c_1}=(E_0^2T_0)^{1/3}(Na^3)^{1/3}L^{2/3}(T_0),$$
$$L(T_0)=\ln\left[\left(\frac{T_0}{E_0}\right)^2\frac{D(T_0)}{Na^3}\right] \quad (1)$$

[where $D(T_0)$ is a logarithmic function of the temperature], the large-scale fluctuations turn out to be suppressed and the state density is determined by the small-scale fluctuations. Here E_0 is the energy of the ground state of the isolated impurity, a is the Bohr radius, and N is the total concentration of the impurities.

The upper limit of the region of the existence of the small-scale fluctuations is determined by the repulsion of the donors from one another and is equal to

$$E_{c_2}=T_0L(T_0). \quad (2)$$

In the region of concentrations that satisfy the inequality $T_0 > E_0(Na^3)$, the condition $E_{c_1} < E_{c_2}$ is satisfied and there exists an energy interval in which the state density is determined by the small-scale fluctuations.

The degree of compensation k needed for the Fermi level to lie in the same interval $E_{c_1} < \mu < E_{c_2}$ can be obtained from the condition that the number of occupied states below the Fermi level be equal to the total number of the electrons (see Ref. 1, p. 373). In the case of a weakly doped semiconductor this condition yields

$$1-k=(Na^3)^{-1}(\mu/E_0)^{3/4}(L(\mu))^{-1/2}\exp\{-(\mu/E_0)^{1/2}L(\mu)\}. \quad (3)$$

It is shown here that the frequency dependence of the absorption coefficient of the electromagnetic radiation is determined by the shape distribution of the small-scale fluctuations having close values of the charge. The absorption line width is small compared with the level spacing. Accordingly, the absorption spectrum has a hydrogenlike structure. The shape of the absorption line corresponding to the $1s \to 2p$ transition is close to Gaussian near the line center and is asymmetric on the line wings.

1. CONNECTION BETWEEN ABSORPTION COEFFICIENT AND THE LEVEL-PAIR DISTRIBUTION FUNCTION

The coefficient of light absorption due to electron transitions from the ground state to the first excited state, within the limits of one small-scale fluctuation, is connected with the real part of the electric conductivity of the crystal by the relation

$$\eta=\frac{2\pi}{n\omega}\operatorname{Re}\sigma(\omega), \quad (4)$$

where

$$\operatorname{Re}\sigma(\omega)=\frac{\pi}{\omega}\iint dE_1\,dE_2\,\rho(E_1,E_2)\,\delta(E_1-E_2-\hbar\omega)$$
$$\times[n_F(E_1)-n_F(E_2)]\,|\langle E_1|j|E_2\rangle|^2. \quad (5)$$

Here $\langle E_1|J|E_2\rangle$ is the matrix element of the current density; $n_F(E)$ is the probability of occupation of an electron level with energy E; $\rho(E_1, E_2)$ is the distribution function of the pair of levels between which the optical transition takes place. The problem reduces thus to a determination of the function $\rho(E_1, E_2)$, which can be obtained by using the optimal-fluctuation method.[2]

The energy spectrum of the small-scale fluctuation is similar to the spectrum of a multiply charged atom,[1] with the ground state for this fluctuation determined by the spherically symmetrical optimal fluctuation. For a spherical fluctuation the first excited state to which an optical transition from the ground state is possible is a p-type threefold degenerate state. In the calculation of the absorption line shape, however, account must be taken of the possible deviations of the shape of the fluctuation from spherical, which leads to a splitting of the degenerate levels of the excited state.

In the theory of disordered systems it is necessary to take into account in the calculation of the state density, with exponential accuracy, only the contribution of the ground states of the fluctuation wells. In the calculation of the absorption coefficient it is necessary to consider the fluctuations that lead to a maximum splitting of the level of the excited state from its position in the sperical well. Since it is not known beforehand which of the three degenerate levels is most strongly detached, it is necessary, in the calculation of the optimal fluctuation, to consider the entire triplet of excited levels.

2. CALCULATION OF THE OPTIMAL-FLUCTUATION FORMULA

According to Ref. 3, the probability of formation of a fluctuation characterized by a specified deviation of the impurity concentration from the mean value $f(\mathbf{r}) = N(\mathbf{r}) - N$ is

$$W\{f\}=\exp[-\Omega\{f\}], \quad (6)$$

where the change in the entropy Ω, due to the onset of the fluctuation f, is equal to

$$\Omega\{f\}=-\int d^3r\left\{[N+f]\ln\frac{N}{N+f}+f\right\}. \quad (7)$$

In the energy region of interest to us the state density is determined by the position of the ground-state level in small-scale fluctuations that are spherically symmetric in shape and are characterized by a total charge Z. To find the level-pair distribution function that determines the light-absorption coefficient according to (5) it is necessary to consider, besides the ground state, also the excited states of the fluctuation well in which the optical transition is allowed. At a given distance between the ground and excited states, equal to the light-quantum energy, the shape of the optimal fluctuation can differ significantly from spherical. In this case the procedure of finding the optimal fluctua-

tion is the following.

The positions of the ground and excited levels of the fluctuation well are determined by the Schrödinger equation

$$-\frac{\hbar^2}{2m}\Delta\Psi+V\{f\}\Psi=\lambda\{f\}\Psi, \tag{8}$$

where $W\{f\}$ is the potential energy of the electron in the well produced by the fluctuation f. In the linear screening approximation the potential $V\{f\}$ is equal to

$$V\{f\}=-\frac{e^2}{\varkappa}\int d^3r' f(\mathbf{r}')\frac{1}{|\mathbf{r}-\mathbf{r}'|}\exp\left\{-\frac{|\mathbf{r}-\mathbf{r}'|}{r_s}\right\}. \tag{9}$$

To solve Eq. (8) we use a variational method and choose hydrogenlike trial functions that satisfy the orthogonality condition

$$\Psi_{1s}=\frac{\beta_1^{3/2}}{\pi^{1/2}}e^{-\beta_1 r},\quad \Psi_{2p^i}=\frac{\beta_2^{5/2}}{\pi^{1/2}}re^{-\beta_2 r}y_i(\vartheta,\varphi);$$
$$y_1=\cos\vartheta,\quad y_{2,3}=2^{-1/2}e^{\pm i\varphi}\sin\vartheta. \tag{10}$$

Here β_1 and β_2 are the variational parameters determined from the condition that the energy eigenfunctions λ_1 and $\lambda_2^{(i)}$ of Eq. (8) be minimal. The level corresponding to the first excited state is split in the field of the nonspherical fluctuation. The optimal fluctuation should minimize the functional $\Omega\{f\}$ of (7) under two conditions:

$$\lambda_1\{f\}=-E_1, \tag{11}$$
$$\lambda_2\{f\}=-E_2, \tag{12}$$

where E_1 and E_2 are fixed positions of the levels of the ground and excited states.

It is not clear beforehand, however, which of the values $\lambda_{1,2}^{(i)}$ corresponds to the largest fluctuation probability. It is therefore necessary to consider the triplet of excited levels simultaneously, and substitute in (12) that value of $\lambda_2^{(i)}$ which corresponds to the highest probability W from (6).

The optimal fluctuation $\tilde{f}(\mathbf{r})$ is determined from the condition of vanishing of the variation

$$\delta\left(\Omega\{f\}+\alpha_1\lambda_1\{f\}+\sum_{i=1}^{3}\alpha_2^{(i)}\lambda_2^{(i)}\{f\}\right)=0, \tag{13}$$

where α_1 and $\alpha_2^{(i)}$ are indeterminate Lagrange multipliers that are subsequently determined from (11) and (12).

Varying (7) and (8) with account taken of (9) and substituting the resultant expressions in (13), we obtain an equation for the optimal fluctuation

$$\tilde{f}(\mathbf{r})=N\left[\exp\left\{\chi_1(r)+\sum_{i=1}^{3}\chi_2^{(i)}(\mathbf{r})\right\}-1\right], \tag{14}$$

where

$$\chi_1(r)=\alpha_1\frac{e^2}{\varkappa}\int d^3r'|\Psi_1(|\mathbf{r}'|)|^2\frac{\exp\{-|\mathbf{r}-\mathbf{r}'|/r_s\}}{|\mathbf{r}-\mathbf{r}'|}, \tag{15}$$

$$\chi_2^{(i)}(\mathbf{r})=\alpha_2^{(i)}\frac{e^2}{\varkappa}\int d^3r'|\Psi_{2^i}(\mathbf{r}')|^2\frac{\exp\{-|\mathbf{r}-\mathbf{r}'|/r_s\}}{|\mathbf{r}-\mathbf{r}'|}. \tag{16}$$

We must thus calculate the functions χ_1 and $\chi_2^{(i)}$. The level-pair distribution function is then obtained by substituting $\tilde{f}(\mathbf{r})$ in Eq. (7).

According to (14)–(16) the difference between the shape of the optical fluctuation and a sphere is determined by the angular dependence of the quantity $\chi_2^{(i)}(\mathbf{r})$. The asymmetry of the fluctuation depends on the ratio of the parameters α_1 and $\alpha_2^{(i)}$. In the calculation of (15) and (16) we shall employ, as usual, the assumption that the optimal fluctuation falls off very rapidly at short distances

$$r<(1/\beta_1)\ll r_s, \tag{17}$$

an assumption that will be justified later on. Substituting (10) in (15) and (16) and calculating the corresponding integrals, we have

$$\chi_1=\frac{\alpha_1 e^2}{\varkappa}\beta_1\left[1-\frac{2}{3}r^2\beta_1^2\right], \tag{18}$$

$$\chi_2^{(1)}(\mathbf{r})=\frac{\alpha_2^{(1)}e^2}{\varkappa}\frac{\beta_2}{2}\left[1+\frac{4}{15}P_2(\cos\vartheta)r^2\beta_2^2\right] \tag{19}$$

$$\chi_2^{(2)}(\mathbf{r})=\frac{\alpha_2^{(2)}e^2}{\varkappa}\frac{\beta_2}{2}\left[1-\frac{2}{15}P_2(\cos\vartheta)r^2\beta_2^2\right], \tag{20}$$

$$\chi_2^{(3)}(\mathbf{r})=\frac{\alpha_2^{(3)}e^2}{\varkappa}\frac{\beta_2}{2}\left[1-\frac{2}{15}P_2(\cos\vartheta)r^2\beta_2^2\right], \tag{21}$$

where $P_2(x)=(3x^2-1)/2$ is the second Legendre polynomial, and $\chi_1(\mathbf{r})$ varies as a function of r over distances $x=(\alpha_1 e^2\beta_1^3/\varkappa)^{-1/2}$. If this characteristic length is to remain less than $1/\beta_1$, we must assume

$$t_1=\alpha_1 e^2\beta_1/\varkappa\gg 1. \tag{22}$$

Substituting (18)–(21) in (14) and using (22), we get

$$\tilde{f}(\mathbf{r})=Ne^{t_1+t_2}\exp\left\{-\frac{r^2}{x^2}[1+2\gamma P_2(\cos\vartheta)]\right\}, \tag{23}$$

where

$$t_2=[\alpha_2^{(1)}+\alpha_2^{(2)}+\alpha_2^{(3)}]\frac{e^2}{\varkappa}\frac{\beta_2}{2}, \tag{24}$$

and the paramter γ characterizes the asymmetry of the fluctuation shape and is equal to

$$\gamma=\frac{1}{5}\left(\frac{\beta_2}{\beta_1}\right)^3\frac{1}{\alpha_1}\left[\alpha_2^{(1)}-\left(\frac{\alpha_2^{(2)}+\alpha_2^{(3)}}{2}\right)\right]. \tag{25}$$

The fluctuation size, which characterizes the shape asymmetry, must also satisfy the condition (17). This imposes the following limitations on the range of variation of the parameter γ:

$$-1+3/2t_1<\gamma<1/2-3/4t_1. \tag{26}$$

3. DISTRIBUTION FUNCTION OF LEVEL PAIRS

We calculate now the energy of the ground state from Eq. (8), using for the potential $V\{f\}$ the expression (9) in which f is replaced by the optimal value of $\tilde{f}$ from (23). We have

$$\lambda_1=E_0a^2\beta_1^2-2ZE_0a\beta_1+2ZE_0a\beta_1^3x^2\frac{1-\gamma}{(1+\gamma)(1-2\gamma)}=-E_1, \tag{27}$$

where Z is the charge of the optimal fluctuation, equal to

$$Z=\int d^3r\,\tilde{f}(\mathbf{r})=Ne^{t_1+t_2}x^3\frac{\pi^{3/2}}{(1+\gamma)(1-2\gamma)^{1/2}}. \tag{28}$$

The variational parameter β_1 is determined from the condition $\partial E_1/\partial\beta_1=0$. It equals, when (22) is taken into

account, the reciprocal of the Bohr radius of a hydrogenlike point atom with charge Z:

$$\beta_1=Z/a. \tag{29}$$

Substituting (29) in (27) we get

$$E_1=E_0Z^2\left[1-\frac{3}{t_1}\frac{1-\gamma}{(1+\gamma)(1-2\gamma)}\right]. \tag{30}$$

The energy of the excited state is calculated by an analogous procedure. For each wave function of the excited state (10) we obtain from the Schrödinger equation (8) the values of $\lambda_2^{(i)}$ and the variational parameter β_2. At our accuracy we have

$$\beta_2=Z/2a \tag{31}$$

and

$$-\lambda_2^{(1)}=\frac{Z^2E_0}{4}\left[1+\frac{3}{20t_1}\frac{\gamma}{(1+\gamma)(1-2\gamma)}\right], \tag{32}$$

$$-\lambda_2^{(2)}=-\lambda_2^{(3)}=\frac{Z^2E_0}{4}\left[1-\frac{3}{40t_1}\frac{\gamma}{(1+\gamma)(1-2\gamma)}\right]. \tag{33}$$

At fixed γ, the fluctuation, in accord with (22), is an allipsoid of revolution, prolate at $\gamma>0$ and oblate at $\gamma<0$ along the z axis. In such a fluctuation the level of the first excited state is split in accord with (32) and (33). At $\gamma>0$ the level $\lambda_2^{(1)}$ shifts downward by double the upward shift of the levels $\lambda_2^{(2)}$ and $\lambda_2^{(3)}$ relative to the positions of the levels for a spherical fluctuation. At $\gamma<0$ the sign of the shift of all three levels is reversed. The largest contribution to the probability (6) of interest to us is made by the level $\lambda_2^{(1)}$, whose shift is maximal. It is therefore necessary in fact to substitute in the condition (11) precisely $\lambda_2^{(1)}$. The equation

$$-\lambda_2^{(1)}=E_2=\frac{Z^2E_0}{4}\left[1+\frac{3}{20t_1}\frac{\gamma}{(1+\gamma)(1-2\gamma)}\right] \tag{32'}$$

together with (30) interrelates the parameters t_1, t_2, and γ. The third relation between them is determined by the condition that the probability (6) be a maximum. To calculate the probability we substitute in (6) the expressions (22) and (29). We obtain

$$\Omega\{\tilde{f}\}=\left(\frac{E_1}{E_0}\right)^{1/2}\left[t_1+t_2-1+\frac{3}{2}\frac{t_2}{t_1}+\frac{3\gamma^2}{(1+\gamma)(1-2\gamma)}\frac{t_1+t_2}{t_1}\right]. \tag{34}$$

From the condition that the maximum condition for the probability (34) we have

$$t_2=\frac{3}{2}\frac{1-\gamma}{(1+\gamma)(1-2\gamma)}. \tag{35}$$

The system (29), (32′) and (35) can be solved in the cases of weak and strong anisotropy of the shape of the fluctuation. In the case of weak anisotropy it is possible to expand in these equations in terms of the parameter $\gamma\ll1$ and retain the terms proportional to γ^2. We then obtain

$$\Omega(E_1,E_2)=\left(\frac{E_1}{E_0}\right)^{1/2}\left[L(E_1)+\frac{2}{3}\left(40\frac{\Delta E_2}{E_1}\right)^2L^2(E_1)\right], \tag{36}$$

where

$$\Delta E_2=E_2-{}^1/_4E_1, \tag{37}$$

and $L(E)$ is defined in (1).

The probability distribution for the position of the level E_2 at fixed E_1 has a Gaussian form. The width of the distribution turns out to be

$$(\overline{\Delta E_2^2})^{1/2}=\frac{E_1^{3/4}E_0^{1/4}}{40(^2/_3)^{1/2}L(E_1)}. \tag{38}$$

The characteristic energy ΔE_1 that determines the exponential fall-off of the state density near the level E_1 is, according to Ref. 1,

$$\Delta E_1=E_0/L^2(E_1). \tag{39}$$

Therefore the width of the distribution of the excited state in terms of the parameter

$$(E_1/E_0)^{1/4}L(E_1)\gg1 \tag{40}$$

is larger than the value of ΔE_1 determined by the ground state. It must be borne in mind, however, that the large numerical factor in (38) causes the width of the distribution of the excited level actually to be larger than ΔE_1 only at very large values of E_1.

In the case of strong anisotropy, the behavior of the function (E_1, E_2) from (34) turns out to be different at $\Delta E_2>0(\gamma\to1/2)$ and $\Delta E_2<0(\gamma\to-1)$. In these two cases,

$$\Omega(E_1,E_2)=\left(\frac{E_1}{E_0}\right)^{1/2}\left[L(E_1)+40\frac{\Delta E_2}{E_1}L(E_1)\right],\quad \Delta E_2>0, \tag{41}$$

$$\Omega(E_1,E_2)=\left(\frac{E_1}{E_0}\right)^{1/2}\left[L(E_1)+80\frac{-\Delta E_2}{E_1}L(E_1)\right],\quad \Delta E_2<0. \tag{42}$$

According to (41) and (42), the distribution of the excited states turns out to be asymmetric at large ΔE_2. We note that just as in the region of small ΔE_2 the distribution of the excited levels depends on the position E_1 of the ground state in the fluctuation.

4. DISCUSSION OF RESULTS

Substituting (36) in (4) we get

$$\mathrm{Re}\,\sigma(\omega)=\frac{\pi}{\omega}\iint dE_1\,dE_2\,\delta(E_1-E_2-\hbar\omega)\,|\langle E_1|j|E_2\rangle|^2$$
$$\times[n_F(E_1)-n_F(E_2)]\exp\left\{-\left(\frac{E_1}{E_0}\right)^{1/2}\left[L(E_1)+\frac{2}{3}\left(40\frac{\Delta E_2}{E_1}\right)^2L^2(E_1)\right]\right\}. \tag{43}$$

Using the expression for the wave functions (10), we can calculate the matrix element $\langle E_1|j|E_2\rangle$, which, naturally, has no exponential dependence on the energy. In our calculations, performed with exponential accuracy, it can be regarded as constant. We then have at $T=0$

$$\mathrm{Re}\,\sigma(\omega)=\frac{\pi}{\omega}|\langle j\rangle|^2\int_{\mu}^{\mu+\hbar\omega}\exp\left\{-\left(\frac{E_1}{E_0}\right)^{1/2}L(E_1)\right\}$$
$$\times\exp\left\{-\frac{2}{3}\left(40\frac{^3/_4E_1-\hbar\omega}{E_1}\right)^2L^2(E_1)\right\}dE_1, \tag{44}$$
$$\langle j\rangle\approx\langle\mu|j|\mu-\hbar\omega\rangle.$$

It is seen from (44) that at the frequency $\hbar\omega_{01}=3/4E_1$ corresponding to the transition between the ground and excited states of the spherical fluctuation, the absorption coefficient is maximal. If the quantity $(\Delta E_2^2)^{1/2}$ from (38) is larger than ΔE_1 from (39), then the quantity that decreases rapidly in the integral (44) is the first factor, which takes into account the fall-off of the state density. Then

$$\operatorname{Re}\sigma(\omega)=\frac{\pi}{\omega}|\langle j\rangle|^2\exp\left\{-\left(\frac{\mu}{E_0}\right)^{1/2}L(\mu)\right\}\cdot$$

$$\times\exp\left\{-\frac{2}{3}\left(40\frac{{}^3/_4\mu-\hbar\omega}{\mu}\right)^2L^2(\mu)\right\}. \quad (45)$$

At $\Delta E_1>(\overline{\Delta E_2^2})^{1/2}$ and at frequencies lower than ω_{01} the line shape is also described by (45). For frequencies higher than ω_{01} the drastically varying quantity is the last exponential factor in (44), which takes into account the broadening due to the deviation of the shape of the fluctuation from spherical. The absorption line shape duplicates in this case the variation of the state density

$$\operatorname{Re}\sigma(\omega)=\frac{\pi}{\omega}|\langle j\rangle|^2\exp\left\{-\left(\frac{4\hbar\omega}{3E_0}\right)^{1/2}L\left(\frac{4}{3}\hbar\omega\right)\right\}. \quad (46)$$

In a region far from the center of the absorption line, at $E_2>1/4E_1(\omega<\omega_{01})$ we have

$$\operatorname{Re}\sigma(\omega)=\frac{\pi}{\omega}|\langle j\rangle|^2\exp\left\{-\left(\frac{\mu}{E_0}\right)^{1/2}L(\mu)\right\}$$

$$\times\exp\left\{-40\frac{[{}^3/_4\mu-\hbar\omega]}{\mu}L(\mu)\right\}, \quad (47)$$

and at $\omega>\omega_{01}$

$$\operatorname{Re}\sigma(\omega)=\frac{\pi}{\omega}|\langle j\rangle|^2\exp\left\{-\left(\frac{\mu}{E_0}\right)^{1/2}L(\mu)\right\}\exp\left\{-80\frac{[\hbar\omega-{}^3/_4\mu]}{\mu}L(\mu)\right\}. \quad (48)$$

It follows from these equations that the optical-absorption line shape should be asymmetrical on the wings and have a slower decrease on the long-wave side. The line width is small compared with the level spacing in accord with the parameter $(E_1/E_2)^{1/4}L(E_1)\gg 1$. For this reason we can assume that higher excited states will also have a small width and will appear in the absorption spectrum in the form of a hydrogenlike series.

For experimental observation of the absorption spectra connected with transitions inside the small-scale fluctuations it is necessary to satisfy a large number of rather stringent conditions. The formulas obtained are valid when the depth of the ground-state level is much less than the width of the forbidden band E_g. On the other hand, the small-scale fluctuations of the multiply charged atom type determine the state density far enough from the edge of the forbidden band. It follows from (3) that at $Na^3\approx 0.1$ and $T_0/E_0\approx 10$ the Fermi level passes through an energy interval in which the state density is controlled by the small-scale fluctuations at a degree of compensation $1-k<10^{-3}$. The concentration of the small-scale fluctuations, which determines the absorption coefficient, is then of the order of $(1-k)N\approx 10^{-3}N$. Since the coefficient of light absorption by the fluctuation wells is close to the absorption coefficient of shallow impurities having the same concentration, observation of this spectrum calls for measurement apparatus of high sensitivity. In semiconductors whose donor activation energy lies in the submillimeter band, such measurements are feasible by using submillimeter technology.[4]

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Dynamic characteristics of EPR signals of saturating systems

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An EPR investigation of the dynamics of the saturation of spin systems has revealed that, in contrast to the Torrey model, the frequency of the oscillations of the absorption signal does not depend on the detuning from resonance, and that there are no oscillations of the dispersion signals. The results are explained by taking into account the nonequilibrium acoustic waves excited in the sample by the radio-frequency field.

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Direct investigations of the dynamics of saturation of spin systems are carried out by observing, in time scale, the establishment of their quasistationary state under the influence of radio-frequency radiation after a jumplike establishment of resonance conditions. A transition process, frequently called Torrey oscillations or transient nutations, can then be observed.[1-4] This effect is connected with the nutational motion of the magnetic-moment vector, which manifests itself in oscillations of the dispersion (χ') and absorption (χ'')

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