Effect of spin-spin interaction pool on NMR saturation in ferromagnets

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A two-temperature theory of NMR saturation in ferromagnets is proposed, in which that part of the spin-spin interaction which yields approximately the dynamic shift of the resonance frequency is included in the modified Zeeman pool. It is shown that the NMR saturation picture corresponding to the two-temperature theory proposed here differs essentially both from the results of the one-temperature theory that includes a dynamic frequency shift, and from the Provotorov theory.

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It is known that the presence of a hyperfine interaction (HFI) between the electron and nuclear spins leads to a number of singularities of the NMR in magnetically-ordered materials. This circumstance was first pointed out by Gossard and Portis[1] and by Freeman and Watson.[2] In particular, it was established that the static part of the HFI is manifest in the form of a strong ($\sim 10^5$–10^6 Oe) static magnetic field. The NMR frequencies in ferromagnets are unusually large ($\sim 10^8$–10^9 sec^{-1}). As to the fluctuating part of the HFI, it produces, on the one hand, an enhancement of the external radio-frequency field, and on the other hand leads to the appearance of an indirect Suhl-Nakamura interaction between the nuclear spins[3,4] which is responsible for the dynamic frequency shift (DFS)[5] and the width of the NMR line.[3] Since the DFS is proportional to the average value of the z component of the nuclear magnetization, the resonant frequency of the nuclear spin system changes, generally speaking, under the influence of the RF field, and this leads to a number of interesting nonlinear effects. Owing to the large radius of the Suhl-Nakamura interaction, the DFS, as a result of the correlation of the motion of the nuclear spins at a distance on the order of the interaction radius, manifests itself also at sufficiently high temperatures of the nuclear spin system, when the latter is disordered to a considerable degree.

The dynamic frequency shift is the main manifestation of the spin-spin interaction in the dynamics of a nuclear spin system that is subject to the interaction of the alternating field, under conditions when the duration of the RF pulse is much shorter than the spin-spin relaxation time. A possible example is provided by experiments on spin echo in ferromagnets, reported a few years ago in the literature.[6] In the case of a more prolonged action of the RF fields it must be borne in mind that the presence of a Suhl-Nakamura interaction leads to the appearance of a thermodynamic subsystem that is connected with the internal degrees of freedom, and which will take active part, together with the Zeeman subsystem, in the energy exchange between the RF field and the spin system. This question is not discussed in the paper of de Gennes et al.,[5] where the theoretical analysis of the NMR saturation in a ferromagnet is based on the one-temperature model of Bloembergen, Purcell, and Pound (BPP)[7] supplemented by the DFS. It must be borne in mind here that (in contrast to a spin system with dipole-dipole interaction) in this case the DFS is appreciable (as noted above) even in the high-temperature approximation. This circumstance makes the problem of separating the spin-spin interaction (SSI) pool a complicated one even in the limit of high temperatures.

In this paper we propose a two-temperature theory of NMR in ferromagnets and show that the presence of the SSI pool alters significantly the saturation picture corresponding to the one-temperature model.

1. The Hamiltonian of a nuclear spin system with Suhl-Nakamura interaction, placed in a static magnetic field $\mathbf{H}_0$, is given by

$$\mathcal{H}_s=-\omega_n\sum_i I_i^z+\sum_{i\neq j}U_{ij}I_i^+I_j^-, \tag{1}$$

where ω_n is the usual NMR frequency.

After the lapse of the spin-spin interaction time, the spin system in question breaks up into an aggregate of thermodynamic subsystems, each of which can be set in correspondence with a thermodynamic parameter, namely the reciprocal temperature. In the case of extremely high temperatures when the DFS can be neglected, these subsystems are the Zeeman pool and the Suhl-Nakamura interaction pool.[1)] In the general case (in the presence of DFS), the picture changes significantly. The frequency of the homogeneous precession of the nuclear spins no longer coincides with ω_n, as a result of which, in transitions induced by the spin-spin interaction, what is conserved is not the usual Zeeman Hamiltonian but the Zeeman Hamiltonian corresponding to the shifted frequency. Indeed, writing down the spin-spin interaction Hamiltonian in the Fourier representation:

$$\frac{1}{2}N^2\sum_k U_k(I_{-k}^+I_k^-+I_{-k}^-I_k^+),$$

where

$$I_k^\pm=\frac{1}{N}\sum_i I_i^\pm\exp(-i\mathbf{k}\mathbf{r}_i),\qquad U_k=\frac{1}{N^2}\sum_{i\neq j}U_{ij}\exp[-i\mathbf{k}(\mathbf{r}_i-\mathbf{r}_j)],$$

we see that the spin-spin interaction component that gives approximately the frequency shift of the homogeneous precession of the nuclear magnetization takes the form[2)]

$$\tfrac{1}{2}U_0(I^+I^-+I^-I^+).$$

It is expedient to include this part of the spin-spin interaction in the Zeeman energy, writing down the initial Hamiltonian of the nuclear spin system (1) in the form[3)]

$$\mathcal{H}_S=\mathcal{H}_z+\tilde{\mathcal{H}}_{SN}, \tag{2}$$

where

$$\mathcal{H}_z=-\omega_n I_0^z+{}^1/_2U(I^+I^-+I^-I^+);$$

$$\tilde{\mathcal{H}}_{SN}=\lim_{\varepsilon\to 0}\int_{-\infty}^{0} dt\, e^{\varepsilon t}\mathcal{H}_{int}(t), \quad \varepsilon\to 0;$$

$$\mathcal{H}_{int}=\sum_{i\neq j} U_{ij}I_i^+I_j^- - {}^1/_2U_0(I^+I^-+I^-I^+),$$

with

$$\mathcal{H}_{int}(t)=\exp(i\mathcal{H}_z t)\mathcal{H}_{int}\exp(-i\mathcal{H}_z t).$$

Here $\tilde{\mathcal{H}}_{SN}$ is that part of the interaction Hamiltonian $\mathcal{H}_{int}$ which is secular with respect to $\mathcal{H}_Z$. The conservation of this part of the interaction is due, as usual, to the fact that actually it is precisely this part which is responsible for the establishment of internal equilibrium in the various subsystems. The processes due to the nonsecular part of the Suhl-Nakamura interaction lead to a change of the Zeeman energy by an amount on the order of the energy corresponding to the DFS, but since this value exceeds the energy of the secular part $\tilde{\mathcal{H}}_{SN}$ of the Suhl-Nakamura interaction, these processes, being very slow, will be disregarded. Therefore $\mathcal{H}_Z$ and $\tilde{\mathcal{H}}_{SN}$ must indeed be regarded as the thermodynamic subsystems that are produced in the nuclear spin system with the Suhl-Nakamura interaction after the spin-spin interaction time has elapsed, under conditions when, on the one hand, the high-temperature approximation is valid, and on the other hand the DFS is significant.

We consider now a ferromagnetic sample placed in an alternating magnetic field. The total Hamiltonian of the system can be represented in the form

$$\mathcal{H}=\mathcal{H}_z+\tilde{\mathcal{H}}_{SN}+\mathcal{H}_h+\mathcal{H}_L+\mathcal{H}', \tag{3}$$

where $\mathcal{H}_h$ is the Hamiltonian of the RF field, treated quantum-mechanically, $\mathcal{H}_L$ is the lattice Hamiltonian, and $\mathcal{H}'$ is the interaction Hamiltonian and consists of the energy of the interaction of the spin system with the alternating field

$$\mathcal{H}_{Ih}=-{}^1/_2\gamma_I(h^+I^-+h^-I^+)$$

and with the lattice

$$\mathcal{H}_{IL}=\sum_i \mathbf{L}_i\mathbf{I}_i.$$

We confine ourselves henceforth to intermediate saturation, in which the probability of the transitions due to the alternating field is smaller than the probability of the transitions due to the spin-spin interaction. The macroscopic evolution stage in which the system in question is represented as an aggregate of weakly interacting subsystems, namely the modified zeeman (MZ) pool $\mathcal{H}_Z$, the SSI pool $\tilde{\mathcal{H}}_{SN}$, the RF field, and the thermostat (lattice), and which is just the one of interest from the point of view of NMR saturation, can be described within the framework of the quantum-statistical method proposed by Zubarev for the construction of a nonequilibrium statistical operator.[9]

Following this method, we arrive after some calculations to the following system of equations describing the time variation of the reciprocal temperatures β_Z of the MZ pool and β_{SN} of the SSI pool:

$$\frac{d\beta_z}{dt}=-\frac{\pi\omega_1^2}{2}f^{+-}(\Omega-\omega_z)\left(\beta_z+\frac{\Omega-\omega_z}{\omega_n}\beta_{SN}\right)-\frac{\beta_z-\beta_L}{T_{zL}},$$

$$\frac{d\beta_{SN}}{dt}=-\frac{\pi\omega_1^2}{2}f^{+-}(\Omega-\omega_z)\frac{\omega_n(\Omega-\omega_z)}{\omega_{SN}^2}\left(\beta_z+\frac{\Omega-\omega_z}{\omega_n}\beta_{SN}\right)-\frac{\beta_{SN}-\beta_L}{T_{SNL}} \tag{4}$$

where

$$f^{+-}(\omega)=\frac{1}{2\pi}\int_{-\infty}^{\infty} dt\, e^{i\omega t}\frac{\mathrm{Sp}[I^-\exp(i\tilde{\mathcal{H}}_{SN}t)I^+\exp(-i\tilde{\mathcal{H}}_{SN}t)]}{\mathrm{Sp}(I^-I^+)}$$

is the form of the absorption line,

$$\frac{1}{T_{zL}}=\frac{\pi}{2}\int_{-\infty}^{\infty} d\omega\, L^{+-}(\omega)f^{+-}(\omega-\omega_z),$$

$$\frac{1}{T_{SNL}}=\frac{\pi}{2}\int_{-\infty}^{\infty} d\omega\, L^{+-}(\omega)f^{+-}(\omega-\omega_z)\frac{(\omega-\omega_z)^2}{\omega_{SN}^2}+\frac{\pi}{2}\int_{-\infty}^{\infty} d\omega\, L^{zz}(\omega)f_i^{zz}(\omega-\omega_z)\frac{(\omega-\omega_z)^2}{\omega_{SN}^2}$$

are respectively the rates of relaxation of the MZ and SSI pools to the thermostat, $\omega_Z = \omega_n + a\beta_Z$ is the NMR frequency, including the DFS, $a\beta_Z \equiv (^2\!/_3)I(I+1)NU_0\omega_n\beta_Z$, Ω is the frequency of the alternating field, and ω_1/γ_I is the amplitude of the alternating field. Here

$$L^{\alpha\beta}(\omega)=\frac{1}{2\pi}\int_{-\infty}^{\infty} dt\, e^{i\omega t}\,\mathrm{Sp}\,\rho_L L^\alpha L^\beta(t), \quad \alpha,\beta=+,-,z;$$

$$\rho_L=\frac{\exp(-\beta_L\mathcal{H}_L)}{\mathrm{Sp}\exp(-\beta_L\mathcal{H}_L)}, \quad \omega_{SN}^2=\frac{\mathrm{Sp}\,\tilde{\mathcal{H}}_{SN}^2}{\mathrm{Sp}(I^z)^2},$$

$$f^{zz}(\omega)=\frac{1}{2\pi}\int_{-\infty}^{\infty} dt\, e^{i\omega t}\frac{\mathrm{Sp}[I_i^z\exp(i\tilde{\mathcal{H}}_{SN}t)I_i^z\exp(-i\tilde{\mathcal{H}}_{SN}t)]}{\mathrm{Sp}(I_i^z)^2}$$

In the concrete applications that follow we shall assume that

$$f^{+-}(\omega)=\frac{1}{\sqrt{2\pi}\Delta}\exp\left(-\frac{\omega^2}{2\Delta^2}\right)$$

where Δ is the second moment of the shape of the Gaussian line (4).

We note that in the derivation of these equations we used the splitting of the correlation functions by replacing $\hat{I}^Z$ by its mean value, and the DFS was retained only in the difference $\Omega - \omega_Z$.

The system (4) constitutes the initial system of equations describing the NMR saturation in ferromagnets in the two-temperature theory and in not too strong RF fields (the case of intermediate saturation). Formally this system of equations coincides with the Provotorov equations, but in fact there is a great difference between them. Indeed, owing to the dependence of ω_Z on β_Z the system (4) is nonlinear, from which it follows that the picture of the NMR saturation in this case will differ strongly from the corresponding linear system of Provotorov's equations.[4] Naturally, it will also differ qualitatively from the NMR saturation picture corresponding to the one-temperature treatment proposed by de Gennes et al.[5]

We note finally that the system (4) differs from the system obtained by Kochelaev and Nigmatulin[10] in that the equation for the rate of change of the reciprocal temperature of the SSI pool contains the factor $\omega_Z - \Omega$ instead of $\omega_n - \Omega$. This is due to the fact that in the cited paper the DFS-inducing part of the dipole-dipole interaction remains included in the dipole-dipole pool. One should expect this difference to be of no significance in some cases (e.g., in the case of a crystal with cubic symmetry).

Completing the discussion of the general problems connected with the system (4), we note that within the framework of the assumed splitting, the first moment of the shape of the line $f^{+-}(\omega)$ relative to the shifted frequency ω_Z is equal to zero. (This circumstance was taken into account when the spin system was broken up into the thermodynamic subsystems $\mathcal{H}_Z$ and $\tilde{\mathcal{H}}_{SN}$.) However, according to estimates given in the monograph of Turov and Petrov,[11] this dependence is insignificant and will be neglected here.

2. As noted above, the system (4) is nonlinear, so that its investigation is a complicated task. We therefore consider first a case of practical interest, that of a sufficiently strong RF field, in which the influence of the relaxation terms on the behavior of the spin subsystems can first be neglected. The system of equations describing the time variation of β_Z and β_{SN} takes during this stage obviously the form

$$\frac{d\beta_z}{dt}=-\frac{\pi\omega_1^2}{2}f^{+-}(\omega_z-\Omega)\left(\beta_z+\frac{\Omega-\omega_z}{\omega_n}\beta_{SN}\right),$$
$$\frac{d\beta_{SN}}{dt}=-\frac{\pi\omega_1^2}{2}f^{+-}(\omega_z-\Omega)\frac{\omega_n(\Omega-\omega_z)}{\omega_{SN}^2}\left(\beta_z+\frac{\Omega-\omega_z}{\omega_n}\beta_{SN}\right). \qquad (5)$$

It follows from it that during the course of the change of the thermodynamic state of the subsystems, the following condition is satisfied

$$\beta_{SN}(t)-\frac{\omega_n}{\omega_{SN}^2}\left[\Omega-\frac{1}{2}(2\omega_n+a\beta_z(t))\right]\beta_z(t)$$
$$=\beta_{SN}(0)-\frac{\omega_n}{\omega_{SN}^2}\left[\Omega-\frac{1}{2}(2\omega_n+a\beta_z(0))\right]\beta_z(0), \qquad (6)$$

where $\beta_Z(0)$ and $\beta_{SN}(0)$ are respectively the initial values of the reciprocal temperatures of the MZ and SSI pools. It follows further from the same system of equations that the action of a sufficiently strong RF field tends to establish the following relation between the reciprocal temperatures of the MZ and SSI pools

$$\beta_z=\frac{\omega_n-\Omega+a\beta_z}{\omega_n}\beta_{SN}. \qquad (7)$$

The condition (6) means that even upon saturation of a spin system that is in thermal equilibrium with the lattice at the initial instant, $\beta_Z(0)=\beta_{SN}(0)=\beta_L$, a change takes place in the temperature of the SSI pool exactly at the center of the magnetic-resonance line $\Omega=\omega_n+a\beta_L$; this differs qualitatively from the result of the Provotorov theory, in which the dipole-dipole pool does not take part, under analogous conditions, in the energy exchange between the RF field and the spin system. The reason is obviously that saturation changes the resonant frequency, so that detuning from resonance occurs and disturbs the SSI pool.

Using relations (6) and (7) with initial conditions $\beta_Z(0)=\beta_{SN}(0)=\beta_L$ we can show that the action of the alternating fields produces a Zeeman temperature determines by the cubic equation

$$\left(\frac{\beta_z}{\beta_L}-1+\frac{\delta}{a\beta_L}\right)^3+2\frac{\omega_{SN}^2}{(a\beta_L)^2}\left[1-\frac{a\beta_L}{\omega_n}-\frac{1}{2}\left(\frac{\delta}{\omega_{SN}}\right)^2\right]$$
$$\times\left(\frac{\beta_z}{\beta_L}-1+\frac{\delta}{a\beta_L}\right)+2\frac{\omega_{SN}^2}{(a\beta_L)^2}\left(1-\frac{\delta}{a\beta_L}\right)=0,$$

where $\delta=\omega_n+a\beta_L-\Omega$ is the initial detuning from resonance.

For small detunings $|\delta|\lesssim\Delta$, to which we confine ourselves here, this equation has a unique real solution

$$\beta_z\approx\left[1-2^{1/3}\left(\frac{\omega_{SN}}{a\beta_L}\right)^{2/3}-\frac{\delta}{a\beta_L}\right]\beta_L. \qquad (8)$$

The corresponding reciprocal temperature of the SSI pool is given by

$$\beta_{SN}\approx-\frac{1}{2^{1/3}}\frac{\omega_n}{a\beta_L}\left(\frac{a\beta_L}{\omega_{SN}}\right)^{4/3}\left[1-2^{1/3}\left(\frac{\omega_{SN}}{a\beta_L}\right)^{2/3}-\frac{\delta}{a\beta_L}\right]\beta_L. \qquad (9)$$

We note that $\beta_Z\approx\beta_L$, whereas $|\beta_{SN}|\gg\beta_L$. This result differs qualitatively from that of the one-temperature theory, in which the temperature of the Zeeman pool can be made in principle arbitrarily large by increasing the intensity of the RF field.[6)]

The system (5) is still quite complicated and cannot be solved analytically. The results of a numerical integration of this system for the quantities $\beta_Z/\beta_L\equiv x$ and $\beta_{SN}/\beta_L\equiv y$ as functions of the RF field are shown in Figs. 1 and 2 ($\tau=t/T_{ZL}$). It follows from these figures that, unlike the Provotorov case, the relation (7) is established non-exponentially. With increasing alternating field intensity $\gamma=\omega_1^2/\Delta^2$ (at a fixed detuning δ) the time of establishment of relation (7) decreases, while the steady-state values of the reciprocal temperatures themselves (just as in Provotorov's theory) do not depend on the intensity of the saturating field.

We note in conclusion that the behavior of the spin system after the establishment of relation (7) between the temperatures is determined by the relaxation effects and (in the simplest case $\Omega=\omega_n+a\beta_L$) is described by the equation

$$\frac{dX}{d\tau}=\frac{X[X^3-\alpha\varkappa(1-\nu)X+\alpha\varkappa]}{X^3+\varkappa},$$

where

$$X=1-\frac{\beta_z}{\beta_L},\quad \tau=\frac{t}{T_{zL}},\quad \alpha=\left(\frac{\omega_{SN}}{a\beta_L}\right)^2,$$
$$\varkappa=\frac{T_{zL}}{T_{SNL}},\quad \nu=\frac{a\beta_L}{\omega_n}.$$

This equation can obviously be easily integrated. We shall not consider this question in this paper, however.

3. We proceed now to consider the stationary state that is established when an RF field of arbitrary inten-

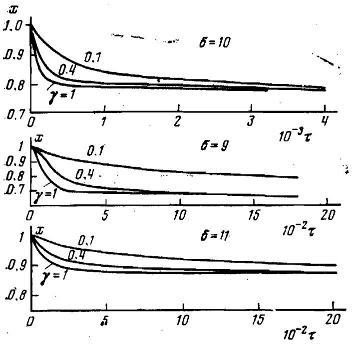

FIG. 1

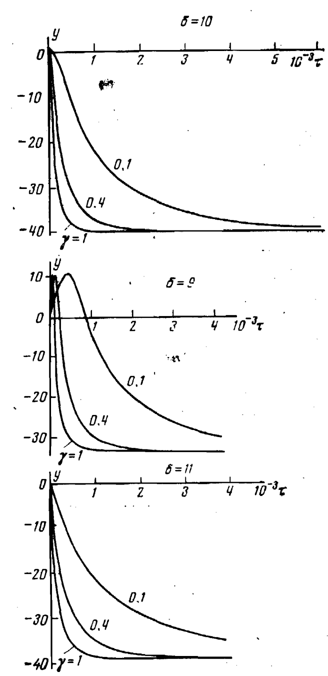

FIG. 2

sity acts on a spin system. Equating to zero the right-hand sides of Eqs. (4) and performing simple calculations, we find that the stationary value of the reciprocal temperature of the SSI pool is connected with the reciprocal temperature of the MZ pool by the relation

$$\beta_{SN}=\beta_L-\frac{T_{SNL}}{T_{zL}}\frac{\omega_n}{\omega_{SN}^2}(\Omega-\omega_n-a\beta_z)\times(\beta_z-\beta_L), \tag{10}$$

whereas the latter is determined by the transcendental equation

$$\sqrt{\frac{\pi}{8}}\frac{\omega_1^2}{\Delta}\exp\left\{-\frac{(\Omega-\omega_n-a\beta_z)^2}{2\Delta^2}\right\}\cdot\left[\beta_z+\frac{T_{SNL}}{T_{zL}}\frac{(\Omega-\omega_n-a\beta_z)^2}{\omega_{SN}^2}(\beta_z-\beta_L)\right]=-\frac{\beta_z-\beta_L}{T_{zL}}. \tag{11}$$

This equation differs from the corresponding equation of the one-temperature theory in the second term of the left-hand side, which determines the effect of the SSI pool on the stationary value of the MZ pool temperature.

It follows from (10) that the temperature of the SSI pool depends to a considerable degree on the ratio of the relaxation times T_{SNL} and T_{ZL}. We shall therefore discuss this question in greater detail.

In ferrodielectrics, the relaxation of the nuclear spins is due mainly to the HFI of the nuclear spins with the magnetic-electron spins. If the quantization axes of the electron and nuclear spins coincide, then the relaxation of the MZ reservoir is determined by the transverse part of the HFI and constitutes a single-magnon process.7) The relaxation of the SSI pool, on the other hand, is due to the longitudinal part of the HFI and is a two-magnon process.8) Owing to the appreciable gap in the spin-wave spectrum, the one-magnon process is slow, so that $T_{SNL} \ll T_{ZL}$. If the quantization axes do not coincide, the ratio of the relaxation time of the MZ and SSI pools, as shown in[11], takes the form $T_{SNL}/T_{ZL} \approx \sin^2\varphi$. Since the angle between the quantization axes is usually small, $\sin^2\varphi \lesssim 0.1$, this ratio remains much less than unity, as before. Consequently, as follows from (10), we have $\beta_{SN} \approx \beta_L$, i.e., the SSI pool is "short-circuited" by the lattice. Under these conditions the contribution of the second term in the right-hand side of (11) is negligible and the latter, naturally, goes over into the equation of the stationary state of the one-temperature theory.[5]

An entirely different situation should take place in ferrometals. Here the relaxation of the nuclear spins is due to their interaction with the conduction electrons (the Korringa mechanism), which practically introduces no distortion in the interaction of the nuclei via the magnetic electrons. In this case $T_{ZL}/T_{SNL} \approx 2$, and, as follows from (11), the SSI pool makes an appreciable contribution to the formation of the stationary state. In particular, single the right-hand side of (11) is positive,9) the stationary state β_Z satisfies the condition

$$\beta_z+\frac{T_{SNL}}{T_{zL}}\frac{(\Omega-\omega_n-a\beta_z)^2}{\omega_{SN}^2}(\beta_z-\beta_L)>0, \tag{12}$$

from which it follows that, regardless of the power of the saturating field, in the two-temperature theory the nuclear spin system with Suhl-Nakamura interaction cannot be fully saturated, i.e., β_Z cannot be made arbitrarily small. This result differs qualitatively from the corresponding result of the one-temperature theory, in which an increase of the intensity of the saturating field can ensure in principle complete saturation of the spin system.

Consider, for example the case when a sufficiently strong field is applied at a frequency $\Omega = \omega_n + a\beta_L$. Under these conditions the stationary value of β_Z is determined approximately by the cubic equation

$$\left(1-\frac{\beta_z}{\beta_L}\right)^3+\frac{T_{zL}}{T_{SNL}}\left(\frac{\omega_{SN}}{a\beta_L}\right)^2\left(1-\frac{a\beta_L}{\omega_n}\right)\left(1-\frac{\beta_z}{\beta_L}\right)-\frac{T_{zL}}{T_{SNL}}\left(\frac{\omega_{SN}}{a\beta_L}\right)^2=0, \tag{13}$$

which has a unique real solution

$$\beta_z\approx\beta_L\left[1-\left(\frac{T_{zL}}{T_{SNL}}\right)^{1/3}\left(\frac{\omega_{SN}}{a\beta_L}\right)^{2/3}\right],$$

which is close in magnitude to β_L. As expected, Eq. (13) coincides with the stationary-state equation for X(t). It can be shown that in the more general case of small detunings ($|\delta| \lesssim \Delta$) from resonance a unique stationary state is also realized.

We note in conclusion that at extremely high temperatures, when the DFS can be neglected, Eq. (13) leads to the result of the Provotorov theory for the stationary value of the Zeeman temperature in the case of intermediate saturation of the magnetic resonance.

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1)In this case the Suhl-Nakamura interaction pool contains the entire energy of the spin-spin interactions.

2)Only homogeneous precession is excited when the alternating field on the nuclear spin system (i.e., in the case of NMR).

3)That the usual separation of the subsystems is incorrect at low temperatures is mentioned in [8].

4)In the case of extremely high temperatures, when the DFS can be neglected, the system (4) coincides with the Provotorov equations.

5)Neglecting the change of the SSI pool temperature, the system (4) reduces to an equation for the Zeeman temperature, which coincides exactly with the equation of the one-temperature theory. [5]

6)The slight difference between β_Z and β_L at small initial detunings from resonance is due to the fact that doring the entire time that relation (7) is established the detuning (which is due to the change of β_Z) does not exceed the line width Δ in order of magnitude.

7)The longitudinal component of the HFI cannot cause transitions with change of the longitudinal magnetization of the nuclei, since it commutes with I^Z.

8)The operator I_i^Z does not commute with $\mathscr{H}_{SN}$.

9)Under the influence of an RF field, β_Z decreases and $\beta_Z - \beta_L < 0$.

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