Anomalous scattering of neutrons in spin-polarized media

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A new exchange mechanism of inelastic scattering with spin flip for slow neutrons propagating through a spin-polarized medium is studied. The scattering is accompanied by emission or absorption of thermal fluctuations of the transverse magnetization of the medium; the weakly damped Larmor precession of nuclear spins in the external magnetic field plays the main role in these fluctuations. Under the conditions of "giant opalescence" the effect is enormous and the corresponding cross sections are significantly greater than the standard elastic scattering cross sections. Thus in the case of ^{29}Si↑ and ^{3}He↑ under typical experimental conditions the cross sections of these inelastic processes are of the order of 10^5–10^6 b.

A new mechanism of transport processes in spin-polarized quantum gases which are accompanied by spin flip was proposed in Ref. 1. This mechanism is based on the emission or absorption of collective spin modes (oscillations of the macroscopic magnetization) by paramagnetic particles forming the gas, but unlike all previous theories it is of a purely exchange origin. The phenomenon under study significantly affects the kinetic properties of the system at quite low temperatures, when the characteristic momenta of the particles are small, $pr_0 \ll \hbar$ (here r_0 is the interaction range). The corresponding cross section of this inelastic process is comparable to and under certain conditions can be many orders of magnitude greater than the typical values of the gas-kinetic elastic scattering cross sections.

In this paper the approach developed previously for the case of quantum gases is employed to describe the properties of a beam of slow neutrons which propagates through a medium with polarized nuclear spins. We will be interested in the change in the density and polarization of the beam as it passes through the target. In other words, our goal is to determine the mean free paths of slow neutrons in a medium with polarized nuclei. Neutrons are knocked out of the beam both by elastic scattering directly by the target nuclei and by interaction with collective thermal fluctuations of the nuclear magnetization in the medium. The contribution of the last mechanism can be gigantic and is observable with modern experimental techniques.

1. The exchange interaction of a neutron and a nucleus (exchange in the sense that the total spin vector of both particles is conserved) plays a very important role. The dependence of the interaction on the total spin of the neutron and the nucleus can be expressed quantitatively in terms of the neutron scattering amplitude $\hat{f}_{si}$ in the following form[2]:

$$\hat{f}_{si}=a+b\hat{\mathbf{s}}\hat{\mathbf{i}},$$
$$a=(2i+1)^{-1}[(i+1)f_+ + if_-],\quad b=2(2i+1)^{-1}(f_+ - f_-). \qquad (1)$$

Here $\hat{\mathbf{i}}$ and $\hat{\mathbf{s}}$ are the nuclear spin and neutron spin operators and $f_\pm$ are the values of the scattering amplitude with total spin equal to $i \pm 1/2$. In the case of slow particles, when the de Broglie wavelength of the neutron $2\pi\hbar/p$ is much larger than the characteristic size of r_0 of the nucleus, the quantities $f_\pm$ are, to a very high degree of accuracy, constants which are independent of the angle and the energy. For definiteness we shall study a target consisting of nuclei with spin $i = 1/2$. In this case the elastic scattering cross sections in the case when the neutron spins are parallel to the nuclear spins as well as in the case when the total spin is zero are given by the obvious formulas

$$\sigma_{\uparrow\uparrow}\equiv\sigma_+=4\pi f_+^2,\quad \sigma_{\uparrow\downarrow}\equiv\sigma_-=4\pi f_-^2,\quad f_\pm\sim r_0. \qquad (2)$$

It was pointed out in Ref. 1 that at low energies an interaction of the spin with the fluctuating magnetic moment exists in the exchange approximation and is very efficient. The exchange Hamiltonian for the interaction of the neutron with the fluctuations of the macroscopic nuclear magnetization can be derived completely analogously to the problem of the interaction of a neutron with the field of density fluctuations. The correction introduced by the quantum refraction of a neutron beam by the nuclei in the target to the characteristic energy of a slow neutron propagating in the medium is given by the expression

$$\hat{H}_s=-(2\pi\hbar^2/\mu)\ \mathrm{Sp}_i\, \hat{f}_{si}\hat{N}_i, \qquad (3)$$

where μ is the reduced mass of the neutron and nucleus and $\hat{N}_i$ is the spinor operator of the density of the medium, which in the general case, in the presence of spin polarization, is a linear function of the spin operator $\hat{\mathbf{i}}$. The density of nuclei N and the magnetic moment $\mathbf{M}$ per unit volume in the target are given by the relations

$$N=\mathrm{Sp}_i\, \hat{N}_i,\quad \mathbf{M}=2\beta\ \mathrm{Sp}_i\, \hat{\mathbf{i}}\hat{N}_i, \qquad (4)$$

where β is the magnetic moment of a nucleus. In the equilibrium state the operator $\hat{N}_i$ is diagonal, as it should be, and has the following form:

$$\hat{N}_i={}^1/_2N(1+2\alpha\hat{\mathbf{i}}\mathfrak{M}),\quad \alpha=(N_+-N_-)/N,\quad N_+ + N_- = N. \qquad (5)$$

Here $\mathfrak{M}$ is a unit vector in the direction of spin polarization and $N_\pm$ are the populations of the states with spins oriented parallel and antiparallel to $\mathfrak{M}$. In the general case the formulas (4) make it possible to express the density operator $\hat{N}_i$ in terms of macroscopic variables

$$\hat{N}_i=N/2+\mathbf{M}\hat{\mathbf{i}}/\beta. \qquad (6)$$

Substituting (1) and (6) into the relation (3) we find finally

$$\hat{H}_s(\mathbf{r},t)=-\frac{2\pi\hbar^2}{\mu}aN(\mathbf{r},t)-\frac{\pi\hbar^2}{\mu\beta}b\hat{\mathbf{s}}\mathbf{M}(\mathbf{r},t). \qquad (7)$$

The first term in (7) is the well-known Hamiltonian for the interaction of a slow neutron with density fluctuations, and the corresponding differential inelastic scattering cross section is expressed in terms of the dynamic structure factor of the medium.[3] The second term in the expression (7) describes the interaction of a neutron with the field of fluctuations of the magnetization. It is obvious that in this case the interaction cross section will be determined by the magnetic form factor of the medium, i.e., by the correlation function of the fluctuations of the magnetic moment. We note that even in the equilibrium case, neglecting the fluctuations of the quantities N and $\mathbf{M}$ terms of the form (7) give rise to some interesting phenomena, such as quantum refraction, rotation of the neutron spin in a polarized target, etc.[2,4]

As will be made evident below, the greatest effect is connected with the inelastic scattering of neutrons by fluctuations of the transverse magnetization, which to first order are not coupled with the density fluctuations. For this reason we retain in the Hamiltonian (7) only the purely magnetic term. The probability of a transition of the neutron from the state $|s_z,\mathbf{p}>$ into the state $<s'_z,\mathbf{p}'|$ is determined by the well-known "golden rule" formula from quantum mechanics

$$dw=\frac{1}{\hbar^2}\left|\int\limits_{-\infty}^{\infty}\langle s_z',\mathbf{p}'|\hat{H}_s|s_z,\mathbf{p}\rangle e^{-i\omega t}\,dt\right|^2\frac{d^3p'}{(2\pi\hbar)^3},\qquad(8)$$

in which the transition frequency $\hbar\omega$ is given by the relations

$$\varepsilon=p^2/2m-2\beta s_zH,\quad \varepsilon'=p'^2/2m-2\beta s_z'H,\quad \hbar\omega=\varepsilon-\varepsilon',\qquad(9)$$

where H is the external magnetic field and m is the neutron mass. It is easy to see that only the matrix elements corresponding to spin flip are nonzero; this is connected with the structure of the Hamiltonian $\hat{H}_S$ in (8). Indeed, the emission or absorption of a spin mode is accompanied by a change in the magnetization $\Delta M_Z=\pm 2\beta$, which is compensated by flipping of the neutron spin. The further calculations are analogous to those performed in Refs. 1 and 5. After substituting the second term of the Hamiltonian (7) into the formula (8) and averaging over the fluctuations of the magnetic moment, we get the correlation function

$$\langle M_i(\mathbf{r}_1,t_1)M_k(\mathbf{r}_2,t_2)\rangle=S_{ik}(\mathbf{r},t),$$

where $\mathbf{r}=\mathbf{r}_1-\mathbf{r}_2$ and $t=t_1-t_2$, and its Fourier component $S_{ik}(\mathbf{k},\omega)$—the dynamic magnetic form factor—appears in the final result. Normalizing the wave function of the neutron in the starting state to unit flux density and dividing the result obtained using the formula (8) by the number of nuclei in the target, we arrive at the differential scattering cross section per nucleus. Thus for the transition $|-1/2,\mathbf{p}\rangle\to\langle 1/2,\mathbf{p}'|$ we obtain with the help of the indicated procedure

$$d\sigma(\downarrow\to\uparrow)\equiv d\sigma_1=\frac{m}{2pN}\left(\frac{\pi\hbar b}{\mu\beta}\right)^2\times[S_{xx}(\mathbf{q},\omega)-iS_{xy}(\mathbf{q},\omega)]\frac{d^3p'}{(2\pi\hbar)^3},\qquad(10)$$

where $\hbar\mathbf{q}=\mathbf{p}-\mathbf{p}'$ is the transferred momentum. The scattering cross section for the process $|1/2,\mathbf{p}\rangle\to\langle-1/2,\mathbf{p}'|$ differs from Eq. (10) only by the sign of the second term

$$d\sigma(\uparrow\to\downarrow)\equiv d\sigma_2=\frac{m}{2pN}\left(\frac{\pi\hbar b}{\mu\beta}\right)^2\times[S_{xx}(\mathbf{q},\omega)+iS_{xy}(\mathbf{q},\omega)]\frac{d^3p'}{(2\pi\hbar)^3}.\qquad(11)$$

Thus finding the angular and energy distributions of the scattered neutrons reduces to calculating the magnetic structure factor of the medium.

Expressing the form factor $S_{ik}(\mathbf{q},\omega)$ in terms of the generalized magnetic susceptibility $\chi_{ik}(\mathbf{q},\omega)$ (see Ref. 1) we obtain

$$d\sigma_1=\frac{m\hbar}{pN}\left(\frac{\pi\hbar b}{\mu\beta}\right)^2(\mathcal{N}_\omega+1)\operatorname{Im}\chi_+(\mathbf{q},\omega)\frac{d^3p'}{(2\pi\hbar)^3}.\qquad(12)$$

Here $\chi_+=\chi_{xx}+i\chi_{xy}$, and the function $\mathcal{N}_\omega$ is determined by the Bose-Einstein formula

$$\mathcal{N}_\omega=(e^{\hbar\omega/T}-1)^{-1}.\qquad(13)$$

If the external magnetic field H is not too weak, so that the uniform Larmor precession of the nuclear spins in the target decays slowly, the susceptibility $\chi_+(\mathbf{q},\omega)$ can be represented in the form

$$\operatorname{Im}\chi_+(\mathbf{q},\omega)=\frac{2\beta^2N\alpha}{\hbar}\frac{\tau^{-1}}{(\omega-\Omega_H)^2+\tau^{-2}},\qquad(14)$$

where $\Omega_H=2\beta H/\hbar$ is the nuclear Larmor frequency and τ is the relaxation time of the transverse magnetization. In an unbounded sample the value of τ is determined by the dipole-dipole interaction and for this reason can be very large owing to the smallness of the nuclear magnetic moment. The relaxation time τ can also contain terms of the form $(-Dq^2)^{-1}$, where D is the nuclear spin diffusion coefficient. Analogous (quadratic in $\mathbf{q}$) spin-wave corrections ω_q also exist in principle in the real part of the spectrum of transverse spin fluctuations ω. These terms are of exchange origin and can play a very important role in systems with delocalized nuclei, for example, in Fermi liquids like ^{3}He, spin-polarized gaseous H↑, D↑, and ^{3}He↑, etc.[6] In systems with localized nuclear spins, however, the spatial dispersion in ω and τ is usually small because the overlapping of the nuclear wave functions is small. The Larmor gap Ω_H in the spectrum of the spin mode is significantly greater than all nonlocal corrections and the inverse relaxation time satisfies $\Omega_H\gg\omega_q,\tau^{-1}$, at least in the typical experimental situation. It is this case, which is predicated on the validity of the formula (14), that we shall have in mind below.

The expressions (12–14) determine the differential cross section for scattering of a slow neutron by thermal fluctuations of the transverse nuclear magnetization

$$d\sigma_1=m\alpha\left(\frac{b}{\mu}\right)^2\left[1-\exp\left(-\frac{\Delta\varepsilon+\hbar\Omega_H}{T}\right)\right]^{-1}\times\frac{\hbar\tau^{-1}}{(\Delta\varepsilon)^2+\hbar^2\tau^{-2}}\frac{p'^2}{p}dp'\frac{do'}{4\pi},\qquad(15)$$

$$\Delta\varepsilon=(p^2-p'^2)/2m.$$

Integrating (15), using the relation

$$\lim_{\delta\to0}\frac{\delta}{x^2+\delta^2}=\pi\delta(x),\qquad(16)$$

we obtain the final expression for the total cross section (per target nucleus):

$$\sigma_1=\pi b^2\left(\frac{m}{\mu}\right)^2\alpha(\mathcal{N}_{\Omega_H}+1)$$

$$=\pi(f_+-f_-)^2\left(1+\frac{m}{M}\right)^2\frac{\alpha}{1-\exp(-\hbar\Omega_H/T)} \qquad (17)$$

where M is the mass of the nucleus. The cross section σ_1 given by Eq. (17) characterizes the exchange process of emission of collective spin modes in the medium by a slow neutron moving straight through it. Such fluctuations of the nuclear magnetization are basically the uniform Larmor procession in an external magnetic field. To avoid misunderstandings, however, we stress that it would be wrong to regard expression (17) as the cross section for the emission of only a quantum of uniform precession by the neutron. The cross section (17) corresponds to the emission of an entire spectrum of spin modes with all possible wave vectors q [in Eq. (15) the integration was performed over all scattering angles]. In talking about the emission of primarily a quantum of uniform precession we have in mind only that the spatial dispersion is small compared with the contribution of the Larmor gap Ω_H in the spectrum of transverse spin fluctuations. In this sense the momentum $\hbar q$ transferred to the spin mode in the inelastic scattering process is assumed to be small, though it can be comparable to the neutron momentum, which should also be small by virtue of the criterion that the particles are slow, $pr_0 \gg \hbar$. For this reason the quantity (17) should be regarded as the leading term in the expansion of the neutron scattering cross section in powers of the small momentum p (and hence q also). Analogously for elastic scattering of slow particles as $p \to 0$ the cross section of the reaction approaches a finite and constant limit—the so-called s-scattering.[2]

In the case of inelastic neutron scattering, accompanied by absorption of a spin mode, we have with the help of (11)

$$d\sigma_2=-\frac{m\hbar}{pN}\left(\frac{\pi\hbar b}{\mu\beta}\right)^2(\mathcal{N}_\omega+1)\,\mathrm{Im}\,\chi_+(-\mathbf{q},-\omega)\,\frac{d^3p'}{(2\pi\hbar)^3},$$

$$\hbar\omega=\Delta\varepsilon-\hbar\Omega_H. \qquad (18)$$

Substituting the expression (14) in Eq. (18) for $(-d\sigma_2)$, we recover expression (15) in which only the sign of the term $\hbar\Omega_H$ in the exponent must be changed, i.e.,

$$d\sigma_2(\Omega_H)=-d\sigma_1(-\Omega_H). \qquad (19)$$

As a result we immediately obtain for the total cross section σ_2

$$\sigma_2=\pi b^2\left(\frac{m}{\mu}\right)^2\alpha\mathcal{N}_{\Omega_H}=\pi(f_+-f_-)^2\left(1+\frac{m}{M}\right)^2\frac{\alpha}{e^{\hbar\Omega_H/T}-1}. \qquad (20)$$

In the thermodynamically equilibrium situation, the target nuclei are polarized with the help of an external magnetic field H, the degree of magnetization α can be expressed with good accuracy as

$$\alpha=\tanh(\hbar\Omega_H/2T). \qquad (21)$$

For this reason at high temperatures $\beta H \ll T$ in the leading-order approximation the cross sections σ_1 and σ_2 are equal to one another

$$\sigma_1=\sigma_2={}^1/_2\pi(f_+-f_-)^2(1+m/M)^2 \qquad (22)$$

and are of the same order of magnitude as the usual elastic scattering cross sections σ_+ and σ_- given by the formulas (2). This means that even under typical experimental conditions the mechanism under study makes the same contribution to the scattering and depolarization of the neutron beam in the target as do the traditional elastic processes, and this effect cannot be neglected. At low temperatures $T \ll \beta H$ the probability of absorption of spin fluctuations σ_2 is exponentially small, while the cross section σ_1 is two times larger than the value given by the formula (22) in the limit $\beta H \ll T$.

The efficiency of the new scattering mechanism is significantly higher under quasiequilibrium conditions when the degree of polarization α of the system is fixed and the external magnetic field is significantly weaker than its thermodynamically equilibrium value (21) calculated for fixed α. Such long-lived spin-polarized states can be realized over times $\tau_\varepsilon \ll t \ll \tau_s$, where τ_ε is the exchange time over which equilibrium over the energies (or momenta) is established and τ_S is the relativistic relaxation time of the longitudinal magnetization. In systems with polarized nuclear spins and not too high concentration of nuclei the quantity of τ_S can reach values of several hours or even days (see below). In such states with moderate values of α but in weak magnetic fields the conditions for the existence of the so-called giant opalescence effect are satisfied; this effect can occur even when the relativistic interaction between a paramagnetic particle and the spin fluctuations is weak.[7] In the present case the exchange nature of the starting Hamiltonian enhances this effect by many orders of magnitude. In the high-temperature limit $T \gg \beta H$ with $\alpha = \text{const}$ we obtain from (17) and (20)

$$\sigma_1=\sigma_2=\pi(f_+-f_-)^2(1+m/M)^2\,\alpha T/\hbar\Omega_H, \qquad (23)$$

which under the foregoing conditions is much greater than the elastic cross sections σ_+ and σ_-.

We also call attention to an analogy between the approach developed in this paper and the theory of photon emission by particles. In the present case the quanta of collective spin waves in the medium play the role of photons. Thus, for example, the cross sections (17) and (20) satisfy the Einstein relations for the probabilities of emission and absorption of quanta of uniform precession (neglecting spatial dispersion).

2. We shall study the problem of the passage of a beam of slow neutrons through a polarized target on the basis of a very rough approximation to the mean-free path, i.e., we shall neglect the real structure of the collision integral. We shall assume that every scattered neutron is knocked out of the beam. The neutron beam can initially be polarized to some degree of polarization γ, which is determined by the concentrations of the neutrons with "up" spins $n_\uparrow$ and "down" spins $n_\downarrow$ with the help of the usual relations

$$n_\uparrow-n_\downarrow=n\gamma, \quad n_\uparrow+n_\downarrow=n. \qquad (24)$$

For definiteness we shall assume that the polarization vector of the target is oriented parallel to that of the beam (and to the external magnetic field). Since the beam density is usually low, $n \ll N$, the change in the spin states of the target nuclei, i.e., the change in the numbers $N_\uparrow$ and $N_\downarrow$ as the neutrons pass by, can be neglected (strictly speaking, this requires that the conditions $n_\uparrow \ll N_\uparrow, n_\downarrow \ll N_\downarrow$ be satisfied). Then on the basis of the assumptions made above the propagation of the neutrons in the target is described by the following simple system of linear equations:

$$dn_\uparrow/dx = -n_\uparrow N_\uparrow \sigma_+ - n_\uparrow N_\downarrow \sigma_- - n_\uparrow N \sigma_2,$$
$$dn_\downarrow/dx = -n_\downarrow N_\downarrow \sigma_+ - n_\downarrow N_\uparrow \sigma_- - n_\downarrow N \sigma_1. \qquad (25)$$

The solution of Eqs. (25) is trivial and leads to the following dependence of the density and degree of polarization of the beam on the distance x:

$$n(x) = {}^1/_2 n_0 F_+(x), \quad \gamma(x) = F_-(x)/F_+(x), \quad n_0 \equiv n(0), \qquad (26)$$

$$F_\pm(x) = (1+\gamma_0)\exp\left(-\frac{N}{2}\sigma_{eff}^{(+)}x\right) \pm (1-\gamma_0)\exp\left(-\frac{N}{2}\sigma_{eff}^{(-)}x\right),$$

$$\gamma_0 \equiv \gamma(0),$$

where the effective cross sections $\sigma_{\text{eff}}^{(\pm)}$ are given by the relations

$$\sigma_{eff}^{(+)} = \sigma_+ + \sigma_- + \alpha(\sigma_+ - \sigma_- + 2\tilde{\sigma}_2), \quad \sigma_2 = \alpha\tilde{\sigma}_2,$$
$$\sigma_{eff}^{(-)} = \sigma_+ + \sigma_- - \alpha(\sigma_+ - \sigma_- - 2\tilde{\sigma}_1), \quad \sigma_1 = \alpha\tilde{\sigma}_1. \qquad (27)$$

At distances which are small compared with the effective mean free path, Eqs. (26)–(27) give

$$n(x) = n_0\{1 - {}^1/_4 Nx[\sigma_{eff}^{(+)} + \sigma_{eff}^{(-)} + \gamma_0(\sigma_{eff}^{(+)} - \sigma_{eff}^{(-)})]\},$$
$$\gamma(x) = \gamma_0 - {}^1/_4 N(1-\gamma_0^2)[\sigma_{eff}^{(+)} - \sigma_{eff}^{(-)}]x. \qquad (28)$$

Under the conditions of giant opalescence with $T \gg \beta H$, when $\sigma_1 \approx \sigma_2 \gg \sigma_+, \sigma_-$, the density of the neutron beam decays according to the law

$$n(x) = n_0 e^{-N\sigma_1 x}, \qquad (29)$$

and the degree of polarization in the leading-order approximation is constant $\gamma(x) \approx \gamma_0 = \text{const}$.

We shall study the quantitative scale of the effect for the example of typical systems with polarized nuclei in the situation when manifestation of giant opalescence can be expected. Vlasenko *et al.* studied the photonuclear magnetism of the magnetically dilute spin system of ^{29}Si in silicon single crystals.[8] The concentration of ^{29}Si was equal to only 4.7%, which corresponds to the nuclear-spin density $N = 6.53 \cdot 10^{20}$ cm^{-3}. This gave long spin relaxations times $\tau_s \approx 30$–60 min which are convenient for performing the experiment. The nuclear spins were polarized with unpolarized light from two incandescent lamps at the temperature of liquid nitrogen ($T = 77$ K). Then the polarized sample was placed in liquid helium ($T = 4.2$ K). Under these conditions the temperature of the nuclear spin system was equal to $\theta \approx 10^{-5}$ K. According to the experimental data on NMR the condition $\Omega_H \tau \gg 1$ already holds very well in such a system in magnetic fields of order $H \gtrsim 10$ G. For this reason the cross sections for inelastic scattering of neutrons by such an object, given by Eqs. (7), (20), and (23), are equal to one another:

$$\sigma_1 = \sigma_2 = {}^1/_2\pi(f_+ - f_-)^2(1+m/M)^2 T/\theta \sim 10^5\ \text{b}, \qquad (30)$$

i.e., they have giant values. The neutron mean free path in this case is given by the expression (29) and is equal to $l \approx 10^{-2}$–10^{-1} cm. It is interesting that only 4.7% of all target nuclei (^{29}Si nuclei which are responsible for this specific scattering) determine the neutron opacity of such thin layers of matter. If the temperature T does not significantly affect the quantity θ, i.e., the degree of nuclear magnetization, then it is even more advantageous to perform the experiment at higher temperatures.

Another typical example is spin-polarized ^{3}He$\uparrow$. In this case absorption of neutrons by nuclei can also play a very important role for the singlet state of the pair of colliding particles. The respective terms $(-n_\uparrow N_\downarrow \sigma_{ab})$ and $(-n_\downarrow N_\uparrow \sigma_{ab})$, where $\sigma_{ab} \approx 5.4 \cdot 10^4$ b is the capture cross section, must be added to the right side of the kinetic equations (25) for $n_\uparrow$ and $n_\downarrow$. The solutions of Eqs. (25) can be represented as before in the form (26), except that the additional terms $\delta\sigma_{\text{eff}}^{(\pm)} = \sigma_{ab}(1 \mp \alpha)$ must be introduced into the definition of the effective cross sections (27). As an illustration we shall consider gaseous ^{3}He$\uparrow$, which is polarized by optical pumping.[9] By specially coating the walls of the experimental chamber the spin relaxation time τ_s can be increased up to two and more days.[10] The necessary condition $\Omega_H \tau \gg 1$ in this case is equivalent to the inequality $\Omega_H \gg Dq^2$. Cutting off the wave vector at the inverse mean free path of atoms in the gas we obtain the lower limit of the magnetic field strength: $\Omega_H \gg \nu$, where $\nu \sim Na^2 v_T$ is the gas-kinetic collision frequency (a is the characteristic particle size and v_T is the thermal velocity). Substituting the typical values of the parameters for gaseous ^{3}He$\uparrow$, namely, $N \sim 10^{16}$ cm^{-3}, $T \sim 1$ K, and $\alpha = 0.5$, we obtain an estimate of the minimum possible magnetic field $H \gtrsim 1$–10 G, for which the giant opalescence effect will be strongest. Calculations using the formula (23) for $H = 10$ G gives $\sigma_1 = \sigma_2 \sim 10^5$–$10^6$, which is much greater than even the capture cross section σ_{ab}. The neutron mean free path in the gas in this case is determined by the inelastic cross sections $\sigma_1 = \sigma_2$ and is equal to only $l \sim 10^2$–10^3 cm (in a very low-density system). In stronger magnetic fields an analogous effect can be observed at room temperature.

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