The absorption mechanism discussed here is not limited to n type materials. A somewhat similar mechanism should also be operative in p type specimens according to the interpretation of elastoresistance given by Adams.[7] Here the mechanism involves scattering from one part of the distorted Fermi surface to another part of the same energy surface. One would expect the relaxation time for this mechanism to be considerably shorter than for intervalley scattering.

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Many-Body Problem in Quantum Mechanics and Quantum Statistical Mechanics

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THIS is a progress report on some work[1] concerning the quantum mechanical calculation of the fugacity coefficients b_l (which correspond to the classical cluster integrals) of a Bose, a Fermi, and a Boltzmann gas at low temperatures. A "binary collision expansion" method is developed which allows for the systematic calculation of b_l as expansions in powers of a/λ, where a represents the parameters of the dimensions of length that characterize the low-energy two-body collision and λ is the thermal wavelength. To any power of (a/λ) the calculation of any specific b_l is reduced to a finite number of quadratures. The method, therefore, is the low-temperature counterpart of the high-temperature expansion of b_l.

By going to the limit $T\rightarrow 0$, the binary collision expansion method can also yield the ground-state energy and ground-state wave function in a systematic expansion. It also supplies information concerning the density of energy levels near the ground state.

The method is applied to the case where the interaction is a hard-sphere interaction with diameter a. The particles are assumed to have spin J, where J is taken to be zero for the Bose and Boltzmann gases, and left arbitrary for the Fermi gas. We use the following notations and units: $\hbar=1$, $m=$ mass of particles $=\frac{1}{2}$, $N=$ number of particles, $V=$ volume of box, $\rho=N/V$, $P_F=$ maximum Fermi momentum $=[6\pi^2\rho/(2J+1)]^{\frac{1}{3}}$, and $\lambda=(4\pi/kT)^{\frac{1}{2}}$. The results are tabulated below:

(A) For the Fermi gas, the fugacity expansion is

$$\begin{aligned}\lambda^3 p/kT=\lambda^3\sum_1^\infty b_l z^l= & -(2J+1)g_{\frac{5}{2}}(-z)\\ & -2J(2J+1)[g_{\frac{3}{2}}(-z)]^2(a/\lambda)\\ & -8J^2(2J+1)g_{\frac{1}{2}}(-z)[g_{\frac{3}{2}}(-z)]^2(a/\lambda)^2\\ & +8J(2J+1)F(-z)(a/\lambda)^2+O(a^3/\lambda^3),\end{aligned} \quad (1)$$

where

$$g_n(z)=\sum_{l=1}^\infty l^{-n}z^l, \quad (2)$$

$$F=\sum_{r,s,t=1}^\infty (rst)^{-\frac{1}{2}}(r+s)^{-1}(r+t)^{-1}z^{r+s+t}. \quad (3)$$

(B) For the Bose gas, the fugacity expansion is

$$\begin{aligned}\lambda^3 p/kT= & g_{\frac{5}{2}}(z)-2[g_{\frac{3}{2}}(z)]^2(a/\lambda)\\ & +8g_{\frac{1}{2}}(z)[g_{\frac{3}{2}}(z)]^2(a/\lambda)^2+8F(z)(a/\lambda)^2+O(a^3/\lambda^3).\end{aligned} \quad (4)$$

(C) For the Bose gas, the pressure and density at the transition point are given by

$$\begin{aligned}\lambda^3 p/kT &=1.34-2(2.61)^2(a/\lambda)+O[(a/\lambda)^{\frac{3}{2}}],\\ \lambda^3\rho &=2.61-4(2.61\pi)^{\frac{1}{2}}(a/\lambda)^{\frac{1}{2}}+O[a/\lambda].\end{aligned} \quad (5)$$

To obtain this expression it was necessary to sum the dominant terms in the fugacity expansion to all orders of (a/λ) near $z=1$.

(D) The ground-state energy per particle for a Fermi gas at a finite density ρ and infinite volume is given by

$$\begin{aligned}E/N= & (3P_F{}^2/5)+8\pi a\rho J(2J+1)^{-1}\\ & \times[1+6(11-2\log_e 2)P_F a/35\pi+O(P_F{}^2a^2)].\end{aligned} \quad (6)$$

(E) The ground-state energy per particle for a Boltzmann gas and for a Bose gas at a finite density and infinite volume is

$$E/N=4\pi a\rho[1+128(\rho a^3)^{\frac{1}{2}}/15\pi^{\frac{1}{2}}+O(\rho a^3)]. \quad (7)$$

The parameters of expansion in Eqs. (6) and (7) are determinable by a simple argument without explicit calculation.

The ground-state wave function and the thermodynamical behavior in cases (D) and (E) near $T=0$ were also obtainable in these computations. Details of the binary collision expansion method and the above calculations will appear in a later publication.

The binary collision expansion method is being applied to a more realistic interaction. Calculation with the Lennard-Jones potential is feasible for b_3 at low T. Work in this direction is under contemplation.

Equations (4) and (6) have been obtained before[2] by the method of pseudopotentials. By using the same method it is not difficult to obtain also Eq. (1). It was emphasized in reference 1 that the method of pseudopotentials is not applicable to all orders of a. With the binary collision expansion method the full range of applicability of the pseudopotential becomes clear. One concludes that it should be possible to obtain also Eqs.

(5) and (7) by the pseudopotential method. This problem was looked into in collaboration with Dr. Kerson Huang. It turns out that it is extremely simple to obtain (7) by the pseudopotential method. A detailed description of this computation will be published shortly.

It should be emphasized that the expansions quoted above are probably all asymptotic expansions. One is led to this conclusion by the following argument. A small and negative value of a corresponds to the case where the force is purely attractive with the scattering length $-a$, for which the gas (for any statistics) collapses. The formula should therefore become meaningless for negative a. This would result if, for example, the physical quantities contain such terms as $\exp.[-a^{-3}\rho^{-1}]$.

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Regeneration of θ_1^0 Mesons by a Magnetic Field*

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THIS letter is to point out a surprising phenomenon that should occur if the θ^0 particle has a spin and a magnetic moment. The sequence of events is shown schematically as follows:

$$\theta^0(\uparrow\uparrow)=\frac{\theta_1(\uparrow\uparrow)+i\theta_2(\uparrow\uparrow)}{\sqrt{2}}\xrightarrow[\text{decay}]{}\theta_2(\uparrow\uparrow)=\frac{\theta^0(\uparrow\uparrow)-\bar{\theta}^0(\uparrow\downarrow)}{i\sqrt{2}}$$

$$\xrightarrow[\text{mag. field}]{}\frac{\theta^0+\bar{\theta}^0}{i\sqrt{2}}=-i\theta_1\xrightarrow[\text{decay}]{}2\pi.$$

A θ^0 meson is made, in associated production with a hyperon, and let us say the spin (first arrow) is up. After $\sim 10^{-10}$ sec, θ_1 decay has left only the θ_2 component of the original wave.[1] This also has spin up. The θ_2 is part θ^0, part $\bar{\theta}^0$, each of which has, of course, spin up. But the magnetic moment (second arrow) of the $\bar{\theta}^0$ will be oppositely directed from that of the θ^0, since it is the antiparticle of the θ^0. Therefore, if the θ_2 beam, in vacuum, encounters a uniform vertical magnetic field, there will be an energy difference $\Delta\epsilon=2\mu H$ between the θ^0, $\bar{\theta}^0$, and a difference $\Delta\epsilon/\hbar$ in their De Broglie frequencies. The θ^0, $\bar{\theta}^0$ components will therefore, in time, get out of phase. Once they are out of phase, the particle is no longer entirely a θ_2, but has a θ_1 component. The θ_1 component then quickly decays into two π mesons.

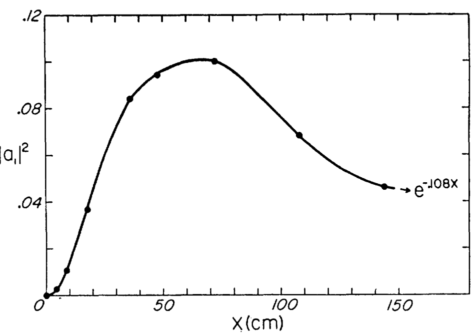

FIG. 1. Regeneration of θ's by magnetic field for $\gamma=6$, $\tau=10^{-10}$ sec, $\omega_2-\omega_1=1/\tau_1=\Delta E_{\text{mag}}$ (which corresponds to 22.7 kg if $\mu=e\hbar/m_Kc$).

A long-lived θ_2 beam can therefore be "quenched" by a magnetic field, if the θ^0 has a magnetic moment.

For a moment $\mu=e\hbar/m_kc$, the fields and times required are about like those needed to rotate a nuclear spin by ~ 1 radian, or on the order of 10^4 or 10^5 gauss-feet. The effect could therefore be observed experimentally.

The equation for the regenerated θ_1 amplitude can be derived in a manner entirely similar to that of Case.[2] A plot of $|\alpha_1|^2$ *vs* distance into the magnetic field (α_1 is the amplitude of θ_1) is given in Fig. 1. The values of the parameters used are: $\tau_1=10^{-10}$ sec, $\tau_2=\infty$ ($\tau_{1,2}=\theta_{1,2}$ lifetime); $\gamma=1/(1-\beta^2)^{\frac{1}{2}}=6$, $\omega_2-\omega_1=$ (mass difference frequency) $=1/2\tau_1$, and $2\mu H/\hbar=1/2\tau_1$ (for $\mu=e\hbar/m_Kc$, this corresponds to a field of 22.7 kilogauss).

Figure 1 implies that if the magnet is 60 cm long, 10% of the incident θ_2's would decay just beyond the magnet. However, $|\alpha_1|^2(1/\tau\gamma\beta c=1/18\text{ cm})$ is the probability, per cm of path, for decay in the magnet. Taking this into account, one sees that the reason the curve is falling, for magnet length $\gtrsim 100$ cm, is that most of the particles have decayed in the magnet. The effect is therefore quite large.

It is interesting to speculate whether experiments already performed can rule out reasonable values of the magnetic moment. A quick look indicates that probably this is not the case. Experiments where θ^0's are produced in a magnet cloud chamber do not have much θ_2 path length. On the other hand, in the experiment by Lederman *et al.*,[3] a strong sweeping magnet preceded the cloud chamber. This could have the effect of wiping out all but the $m=0$ substate (which is unaffected by the field). This substate would then not regenerate in the cloud chamber magnet, since its field was parallel to that of the sweeping magnet. Rather, the typical θ_2 three-body decays would be seen.

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