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SPACE-ENERGY DISTRIBUTION OF NEUTRONS IN A HEAVY GASEOUS MODERATOR

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THE theory of neutron thermalization in a heavy (atomic weight $M \gg 1$) monoatomic gas with constant mean free path λ and constant neutron lifetime τ has been discussed in a number of papers.[1-3] However, the majority of the results refers to the energy distribution only. The space-energy distribution function has been found only in the region of relatively large energies.[2] In the case of weak absorption this problem can be solved exactly.

The equation for the space-energy distribution function[2] may be written as follows:

$$-\alpha\psi(\mathbf{r}, x) + \beta x \nabla^2 \psi(\mathbf{r}, x) + (3 - 2x^2)\frac{\partial\psi(\mathbf{r}, x)}{\partial x} + x\frac{\partial^2\psi(\mathbf{r}, x)}{\partial x^2} = 0, \qquad \alpha = (2M\lambda/\tau)\sqrt{m/2kT}, \quad \beta = 2M\lambda^2/3. \tag{1}$$

where x^2 is the neutron energy in units of kT (T — temperature of the moderator), $\psi(\mathbf{r}, x)$ is the space-energy distribution function divided by $x^2 e^{-x^2}$, and m is the neutron mass.

For a moderator of finite dimensions one may obtain a solution of Eq. (1) in the form of an expansion in a complete set of orthonormal functions $R_l(\mathbf{r})$ of the Laplacian for the corresponding boundary value problem $[\nabla^2 R_l(\mathbf{r}) + \Omega_l R_l(\mathbf{r}) = 0]$, i.e.,

$$\psi(\mathbf{r}, x) = \sum_l R_l(\mathbf{r})\, n_l(x). \tag{2}$$

then each of the functions $n_l(x)$ should satisfy the equation

$$x d^2 n_l/dx^2 + (3 - 2x^2)\, dn_l/dx - (\alpha + \beta x \Omega_l)\, n_l = 0. \tag{3}$$

Making use of the requirement that $n_l(x)$ be finite as $x \to 0$, this equation may be transformed into an integral equation of the Volterra type

$$n_l(x) = C_l \Phi(a, 2, x^2) + \alpha \int_0^x n_l(t) K(x, t)\, dt,$$
$$K(x, t) = \frac{1}{2}\Gamma(a)\, t^2 e^{-t^2} [\Psi(a, 2, t^2)\, \Phi(a, 2, x^2) - \Phi(a, 2, t^2)\, \Psi(a, 2, x^2)], \tag{4}$$

the solution of which, as is well known, is of the form

$$n_l(x) = C_l \sum_{m=0}^{\infty} \alpha^m \varphi_m(x), \quad \varphi_0(x) = \Phi(a, 2, x^2); \quad \varphi_{m+1}(x) = \int_0^x \varphi_m(t) K(x, t)\, dt. \tag{5}$$

Here $\Phi(a, b, z)$ is confluent hypergeometric function and

$$\Psi(a, b, z) = \frac{1}{\Gamma(a)} \int_0^\infty e^{-zt} t^{a-1} (1 + t)^{b-a-1} dt$$

is the other linearly independent solution of the hypergeometric equation; $a = \beta\Omega_l/4$.

The normalization constants C_l may be obtained by comparing the assymptotic form of $n_l(x)$ for large x with the results of a calculation in the Fermi Age approximation. Thus, neglecting absorption for the sake of simplicity, one has for the case of a unit intensity source of neutrons at the point $\mathbf{r} = \mathbf{r}_0$ in the age approximation, the following

$$\psi_{\text{Age}}(\mathbf{r}, x) = \frac{\lambda M}{2}\sqrt{\frac{m}{2kT}}\frac{x^{s}}{x^{4}}\sum_{l} R_l(\mathbf{r})\left(\frac{x}{x_0}\right)^{\beta\Omega_l/2}, \tag{6}$$

where x_0 is the source neutrons speed. Hence

$$\psi(\mathbf{r}, x) = \frac{\lambda M}{2}\sqrt{\frac{m}{2kT}}\sum_{l} R_l(\mathbf{r}_0)\, R_l(\mathbf{r})\, x_0^{-\beta\Omega_l/2}\Gamma\left(\frac{\beta\Omega_l}{4}\right)\Phi\left(\frac{\beta\Omega_l}{4}, 2, x^2\right) \tag{7}$$

In the case of a source located in an infinite homogeneous medium the sum over l must be replaced by the corresponding integral.

A detailed discussion of applications of the above results to various special cases will be published later.

In conclusion I express deep gratitude to F. L. Shapiro for valuable discussions in the process of this work.

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INFLUENCE OF FINITE NUCLEAR SIZE ON EFFECTS CONNECTED WITH PARITY NONCONSERVATION IN BETA DECAY

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For β-decay, particularly for forbidden transitions, the action of the nuclear field is of importance. As in Ref. 1, we take for the electron wave function

$$\psi_e = \begin{pmatrix}\varphi_e\\ \chi_e\end{pmatrix}, \quad \begin{aligned}\varphi_e &= [\alpha_0 + i\mathbf{pr}\alpha_1 + i(\sigma\mathbf{r})(\sigma\mathbf{n})\beta_c]\,u_\xi,\\ \chi_e &= [\beta_0 + i\mathbf{pr}\beta_1 + i(\sigma\mathbf{r})(\sigma\mathbf{n})\alpha_c](\sigma\mathbf{n})\,u_\xi,\end{aligned} \tag{1}$$

where

$$\mathbf{n} = \frac{\mathbf{p}}{p}, \quad \alpha_0 = \sqrt{\frac{\pi}{2p\varepsilon}}\, ie^{-i\delta_{-1}}g_{-1}, \quad \beta_0 = \sqrt{\frac{\pi}{2p\varepsilon}}\, f_1 e^{-i\delta_1}, \quad \alpha_1 = \sqrt{\frac{\pi}{2p\varepsilon}}\frac{3}{pr}\, e^{-i\delta_{-2}}g_{-2}, \quad \beta_1 = \sqrt{\frac{\pi}{2p\varepsilon}}\frac{3}{ipr}\, e^{-i\delta_2}f_2,$$

$$\beta_c = \sqrt{\frac{\pi}{2p\varepsilon}}\frac{1}{r}\left(e^{-i\delta_1}g_1 - e^{-i\delta_{-2}}g_{-2}\right), \quad \alpha_c = \sqrt{\frac{\pi}{2p\varepsilon}}\frac{1}{ir}\left(e^{-i\delta_{-1}}f_{-1} - e^{-i\delta_2}f_2\right)$$

($m_e = c = \hbar = 1$), g_κ, f_κ are the inside-the-nucleus solutions of the radial Dirac equation joined with the outside solutions at $r = r_0$; δ_κ is the phase.[2] We write the β-interaction Hamiltonian as follows

$$H = \sum\left\{g_i(\bar{\psi}_2 O_i \psi_1)\left(\bar{\psi}_e O_i \frac{1-\gamma_5}{2}\psi_\nu\right) + g_i'(\bar{\psi}_2 O_i \psi_1)\left(\bar{\psi}_e O_i \frac{1+\gamma_5}{2}\psi_\nu\right)\right\} \tag{2}$$

(summation over i = S, T, V, A, P). If the two-component neutrino theory[3-5] holds, then emission of an antineutrino together with an electron corresponds to $g_i' = 0$, whereas emission of a neutrino corresponds to $g_i = 0$.