of 19.5±0.2 Mev.[4] There are no associated low-energy electron or recoil tracks at point A (Fig. 2). The absence of a recoil and low energy electron in conjunction with the mass measurements along the primary make it improbable that the proton was ejected by the nuclear capture of a negative π or K meson. In terms of known decay schemes, it seems reasonable to interpret the event as the decay, from rest, of a charged hyperon[5] (Y^+) into a proton and a neutral π meson:

$$Y^+ \rightarrow p + \pi^0 + Q.$$

The Q value for this reaction is found to be 117±2 Mev. From the measured Q value, the mass of the hyperon is found to be $2329\pm5m_e$. The Q value is in excellent agreement with two other similar events reported by the Genoa-Milan and the Padua groups.[6] In these two events, the proton energies were found to be 18.7±0.2 Mev and 18.5±0.3 Mev, respectively. The consistency in the ranges of the protons among the three events (the Genoa-Milan, the Padua event, and the above event) strongly indicate that these three hyperons decayed from rest and that only two particles were involved in the decay.

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Theory of Coulomb Excitation

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THE theory of Coulomb excitation has been considered earlier by many authors. Especially helpful for the interpretation of the experiments is a semiclassical calculation by Ter-Martirosyan.[1,2] In this theory the trajectory of the impinging particle is described by a classical hyperbolic orbit.

We have obtained a somewhat more accurate description of the process by approximating the Coulomb wave function, entering in the quantum-mechanical treatment, by means of the WBK method. The results of this calculation are very similar to the results of the semiclassical theory. The cross sections for electric dipole and electric quadrupole excitation are:

$$\sigma(E1)=\frac{2\pi^2 Z_1^2 e^2}{9\hbar^2 v_i^2}B(E1)g_{E1}(\xi), \tag{1}$$

$$\sigma(E2)=\frac{2\pi^2 m_1^2 v_f^2}{25 Z_2^2 e^2 \hbar^2}B(E2)g_{E2}(\xi), \tag{2}$$

with the following new definition of the parameter:

$$\xi=\alpha_f-\alpha_i=\frac{Z_1 Z_2 e^2}{\hbar}\left(\frac{1}{v_f}-\frac{1}{v_i}\right), \tag{3}$$

where Z_1e and Z_2e are the charges of the impinging projectile and the nucleus, respectively. The initial and final relative velocities are denoted by v_i and v_f, while m is the reduced mass. $B(E\lambda)$ is the reduced transition probability for electric 2^λ-pole radiation in the notation of Bohr and Mottelson.[3] The functions $g_{E\lambda}(\xi)$ are the same as those entering in the semiclassical expression.[1,2] For small excitation energies ΔE, the formulas (1), (2), and (3) reduce to the corresponding semiclassical formulas. We have then $v_f \sim v_i \sim v$, and

$$\xi=\frac{Z_1 Z_2 e^2}{\hbar v}\frac{\Delta E}{2E}. \tag{4}$$

The expression (1) for electric dipole excitation has been compared with the exact quantum mechanical formula[4]

$$\sigma(E1)=-\frac{128\pi^4}{9}\frac{Z_1^2 e^2}{\hbar^2 v_i^2}B(E1)\left(\frac{\alpha_f}{\alpha_i}\right)^2 \times\frac{e^{2\pi\alpha_i}}{(e^{2\pi\alpha_i}-1)(e^{2\pi\alpha_f}-1)}\frac{d}{dx_0}|F(x_0)|^2,$$

where

$$x_0=-4\alpha_i\alpha_f/(\alpha_i-\alpha_f)^2, \tag{5}$$

and

$$F(x_0)={}_2F_1(-i\alpha_i, -i\alpha_f, 1, x_0)$$

in terms of ${}_2F_1$, the ordinary hypergeometric function. In the range $\alpha_i \geqslant 1$ and $0 \leqslant \alpha_f-\alpha_i=\xi \leqslant 1$, the WBK expression (1) agrees within 2 percent with the quantum mechanical calculation, which represents a considerable improvement over the semiclassical treatment.

In the case of the quadrupole excitation, no exact quantum mechanical treatment has been given, but the adequacy of the WBK treatment for the $E1$ case, suggests that (2) should also represent a good approximation when the effects resulting from penetration of the bombarding particle into the nucleus itself can be neglected.

Indeed, it is found that (2) satisfactorily represents all the experimental yield curves available to us at the present time.

Other kinds of electromagnetic excitations have also been considered and a detailed account of all the calcu-

lations will be published in the proceedings of the Danish Academy.[5,6]

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Rf-Induced Transitions of Nuclear Spins at the Electronic Resonance Frequency*

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A REASONABLE explanation is offered to account for the observed increased transition rate of nuclear spins at the electronic spin resonance frequency, as reported in the preceding Letter by Honig. It is known from the work of Bloembergen[1] that an important mechanism of nuclear spin-lattice relaxation in solids is through the coupling of the nuclear magnetic moments to the electron magnetic moments of paramagnetic impurities. The power spectrum associated with the relaxation of the electron moment determines the rate at which nuclear spins are flipped. If the electron spin-lattice relaxation time is very long, even a weak rf field at the electron resonance frequency will have a major effect on the power spectrum. In this way nuclear flips are enhanced when the electron resonance is observed."

The Hamiltonian for a single arsenic site, when the dipolar interactions are neglected, is

$$\mathcal{H}=a\mathbf{I}\cdot\mathbf{S}+H_0(g_e\mu_B S_z-g_N\mu_N I_z) +H_x e^{i\omega t}g_e\mu_B S_x-H_x e^{i\omega t}g_N\mu_N I_x, \quad (1)$$

where the nuclear moment is considered to be positive. For the experimental magnetic field of 3000 oersteds the use of I_z and S_z as good quantum numbers gives the correct energy levels within 2 to 3 percent. The off-diagonal interaction term $\frac{1}{2}aS^+I^-$ can be made to induce nuclear transitions to first order in the sense of time-dependent perturbation theory, by considering the rf-induced transitions of S^+ as a source of power. This is similar to the argument made by Bloembergen, Purcell, and Pound[2] in which they use the time dependence of the magnetic dipolar interactions caused by the rotational and translational motion of the spins to obtain the power necessary for a nuclear spin relaxation. The correlation function for $S^+(t)$ is assumed, following Bloembergen,[1] to be proportional to

$$|S^+|^2e^{i\omega_0 t}\exp(-t/T_{\rm rf}^{\rm el}),$$

where $T_{\rm rf}^{\rm el}$ is the electronic transition time given as[2]

$$T_{\rm rf}^{\rm el}=(\tfrac{1}{2}\mu_B^2H_1^2T_2^{\rm el}/\hbar^2)^{-1}, \quad (2)$$

and $T_2^{\rm el}$ is the half-maximum value of the electronic line shape function.

Corresponding to this correlation function is the power spectrum,

$$|S^+|^2\frac{2T_{\rm rf}^{\rm el}/\pi}{1+4(T_{\rm rf}^{\rm el})^2(\omega-\omega_0)^2}, \quad (3)$$

where ω_0 is the electronic Larmor frequency. Substituting into the standard formula for transition probability,[3] we have for the nuclear transition time,

$$T_{\rm nuc}=\left[\frac{2\pi}{\hbar^2}\left|\frac{a}{2}S^+I^-\right|^2\frac{2T_{\rm rf}^{\rm el}/\pi}{1+4(T_{\rm rf}^{\rm el})^2(\omega_0'-\omega_0)^2}\right]^{-1}. \quad (4)$$

Note that the component of the power spectrum is taken at the nuclear Larmor frequency ω_0'.

Next we will show how the same result can be obtained by second-order perturbation theory. Here we consider the transition between the initial spin state, $m_S=\frac{1}{2}$, $m_I=\frac{3}{2}$, and the final state, $m_S=\frac{1}{2}$, $m_I=\frac{1}{2}$. If standard perturbation theory is used, the expansion coefficient for the final state satisfies the equation,

$$\dot{a}_f^{(2)}=(i\hbar)^{-1}\sum_n H_{fn}a_n^{(1)}e^{i\omega_{fn}t}, \quad (5)$$

where

$$H=H_1g_e\mu_BS_xe^{i\omega t}+\tfrac{1}{2}aS^+I^-+\tfrac{1}{2}aS^-I^+,$$

and

$$a_n^{(1)}=(i\hbar)^{-1}\Big\{\langle n|H_1g_e\mu_BS_x|i\rangle \times\left[\frac{\exp[i(\omega+\omega_{ni}^{(1)}+i/2T_{\rm rf}^{\rm el})t]-1}{i(\omega+\omega_{ni}^{(1)}+i/2T_{\rm rf}^{\rm el})}\right] +\langle n|\tfrac{1}{2}aS^+I^-+\tfrac{1}{2}aS^-I^+|i\rangle \times\left[\frac{\exp[i(\omega_{ni}^{(2)}+i/2T_{\rm rf}^{\rm el})t]-1}{i(\omega_{ni}^{(2)}+i/2T_{\rm rf}^{\rm el})}\right]\Big\}.$$

It is to be noted that we have represented the initial state by a decaying exponential.[4] The only intermediate state which gives a nonzero contribution is $m_S=-\frac{1}{2}$, $m_I=\frac{3}{2}$, and thus only the matrix element containing H_1 contributes to $a_n^{(1)}$. Therefore we set $\omega_{ni}^{(1)}=-\omega_0$ and $\omega_{fn}=\omega_0-\omega_0'$. The transition probability is wanted at the electronic renance frequency, so setting $\omega=\omega_0$ the expansion coefficient then satisfies the equation

$$\dot{a}_f^{(2)}=-(i\hbar)^{-2}(2T_{\rm rf}^{\rm el})\langle f|\tfrac{1}{2}aS^+I^-|n\rangle\langle n|H_1g_e\mu_BS_x|i\rangle \times[\exp(-t/2T_{\rm rf}^{\rm el}+i\omega_0t-i\omega_0't)-\exp(i\omega_0t-i\omega_0't)]. \quad (6)$$

The first two terms can be made to conserve energy between the initial and final states by choosing the Fourier component of $\exp(-t/2T_{\rm rf}^{\rm el})\exp(i\omega_0t)$ at the frequency ω_0'. Formally carrying this out gives just the transition probability in Eq. (4).