representative for the order of magnitude, and corresponds to the maximum distance between the catwhiskers, emitter and collector, where still a transitor interaction may be found. Incidentally, the high quantum yield observed by S. Benzer[3] at illuminated high back voltage germanium detectors, may result from the fact that there is a large probability for the holes released within R (and not only for those released within the barrier layer) to reach the collector before recombination.

The production of holes near the emitter can be understood by the following considerations. The extent of decrease of electron density toward the electrode, which is characteristic of the barrier layer, is large enough so that the hole density, which is negligible in the bulk germanium, becomes higher than the electron density.[1] Thus the current will be carried mostly by holes in the region very close to the electrode, whereas practically only by electrons, in the bulk germanium at sufficient distance from the electrode. The transition of hole current into electron current requires an excess of the recombination of holes with electrons over the rate of production by thermal agitation in the lattice, if a current passes in the forward direction (holes near the electrode and electrons of the bulk germanium move in direction together to each other). On the other hand, there is an excess of formation of holes and electrons if a current flows in the blocking direction (holes near the electrode and electrons of the bulk germanium move in direction apart from each other). Excess of recombination necessitates a density of holes and electrons higher than in thermal equilibrium. Thus the current in the forward direction shifts the equilibrium of holes and free electrons in the neighborhood of the catwhisker in a similar way as a temperature rise would do.

Finally, it should be mentioned that the density of excitons increases with the densities of holes and free electrons. Thus, holes and electrons may find their way as excitons from the emitter area to the collector base field, and they may contribute to current amplification if they dissociate there.

[1] J. Bardeen and W. H. Brattain, Phys. Rev. **74**, 230 (1948).

[2] Quoted by W. Shockley at the meeting of the Am. Phys. Soc., January 27, 1949, in discussion of his presented paper: "Electronic Theory of the Transistor."

[3] S. Benzer, Phys. Rev. **70**, 105 (1946).

On the Difficulty of the Meson Theory of Nuclear Forces

GENTARO ARAKI
Department of Industrial Chemistry, Kyoto University, Kyoto, Japan

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It has generally been believed that the meson theories of nuclear forces involve a difficulty of r^{-3}. Mixtures of two fields were considered first by Møller and Rosenfeld[1] and later by Schwinger[2] in order to eliminate such a term. Furthermore, a divergence difficulty was noticed by the present author, and another mixture was proposed by him.[3] It will be shown in the following, however, that these difficulties are merely apparent so far as the pseudoscalar theory is concerned. Consequently it seems that they have no connection with the general defect of the contemporary field theory.

In the meson theory the matrix element of the two-nucleon potential, W, is given by

$$W_{BA}=\Sigma_I\, H_{BI}'H_{IA}'/(E_I-E_A)-\Sigma_I\, H_{BI}'H_{IBA}'/(E_I-E_A) \quad (1)$$

where H' is the interaction between mesons and nucleons. We can find W itself by the Fourier theorem as follows:

$$W=\Sigma_B\, W_{BA}\exp(i\mathbf{kx}), \quad (2)$$

where $\mathbf{x}$ is the relative position vector between two nucleons, $\hbar\mathbf{k}$ is the momentum of a meson in the intermediate state, I, and matrix elements of the charge and spin operators are suitably converted to the operators themselves.

If we adopt the pseudoscalar meson theory, H' is given by

$$H'=-(f/\kappa^2)\int\psi^\dagger\tau\boldsymbol{\sigma}\,\mathrm{grad}\,U\psi dV+\text{its adjoint} \quad (3)$$

in the non-relativistic approximation where ψ is the wave function of nucleons, U is the pseudoscalar wave function of mesons, and the other notations have the usual meaning. Matrix elements of H' in (1) should be calculated in the non-relativistic approximation, for W is required only in such an approximation. In this approximation we have

$$(1/\kappa)\boldsymbol{\sigma}\,\mathrm{grad}\,U=(2M/\mu)\rho_2 U, \quad (4)$$

where M and μ are masses of a nucleon and a meson, respectively. This identity can easily be shown for every one of their matrix elements with respect to eigenstates of the nucleon momentum on the basis of Dirac's equation. Such a relation was first suggested by Nelson[4] and was later derived by Dyson[5] according to a method of a transformation.

Making use of the relations (1), (2), (3), and (4) we have

$$W=-(2M/\mu)^2(f^2/r)\exp(-\kappa r)T\rho_2^{(1)}\rho_2^{(2)}, \quad (5)$$

where $T=(\tau_1^{(1)}\tau_1^{(2)}+\tau_2^{(1)}\tau_2^{(2)})/2$ (charged theory) or $\boldsymbol{\tau}^{(1)}\boldsymbol{\tau}^{(2)}/2$ (symmetrical theory). At first sight it seems as if this potential vanishes in the non-relativistic approximation, but this is not the case because of the first factor and of the peculiar properties of ρ_2 and Dirac's Hamiltonian. This two-nucleon potential has neither an r^{-3} term nor a divergent integral (especially it has no term of δ-function type). This verifies what was said at the beginning of this letter.

The potential given by (5) seems to be quite different from a usual one, but it is just the same as a hitherto well-known expression, as will be shown in the following. The Hamiltonian of a two-nucleon system with the interaction (5) is given by

$$H=\rho_1^{(1)}\boldsymbol{\sigma}_1\mathbf{p}_1+\rho_1^{(2)}\boldsymbol{\sigma}_2\mathbf{p}_2+(\rho_3^{(1)}+\rho_3^{(2)})Mc^2+W, \quad (6)$$

where $\mathbf{p}$ is equal to $-i\hbar c\nabla$. We expand the eigenfunction of H as follows:

$$\psi=\psi^{++}\chi^+(1)\chi^+(2)+\psi^{+-}\chi^+(1)\chi^-(2)+\psi^{-+}\chi^-(1)\chi^+(2)+\psi^{--}\chi^-(1)\chi^-(2), \quad (7)$$

where χ^+ and χ^- are eigenfunctions of ρ_3 belonging to its eigenvalues $+1$ and -1, respectively. If we substitute (6) and (7) into $H\psi=E\psi$, this equation becomes a system of simultaneous equations for ψ^{++}, ψ^{+-}, ψ^{-+}, and ψ^{--}. In the non-relativistic case ψ^{++} is largest. If we eliminate ψ^{+-}, ψ^{-+}, and ψ^{--} on an assumption that $\phi\ll 1$, we have $H''\psi^{++}=(E-2Mc^2)\psi^{++}$ in the non-relativistic approximation where H'' is a non-relativistic Hamiltonian given by $H''=-(\hbar^2/2M)\times(\Delta_1+\Delta_2)+W'$. W' and ϕ are given by

$$W'=f^2T\exp(-\kappa r)/(1+\phi)(1+2\phi)r\{\tfrac{1}{3}\boldsymbol{\sigma}_1\boldsymbol{\sigma}_2+(\tfrac{1}{3}+1/\kappa r+1/\kappa^2r^2)\Lambda\}, \quad (8)$$

$$\phi=(M/\mu)(f^2/\hbar c)(|a|/\kappa r)\exp(-\kappa r), \quad (9)$$

and

$$\Lambda=3(\boldsymbol{\sigma}_1\mathbf{x})(\boldsymbol{\sigma}_2\mathbf{x})/r^2-\boldsymbol{\sigma}_1\boldsymbol{\sigma}_2, \quad (10)$$

where a is the eigenvalue of T.

ϕ in the expression of W' should be omitted, for W' is derived in the non-relativistic approximation. If we put $\phi=0$ in Eq. (8) it is just the same as the well-known potential of the pseudoscalar meson theory. Equation (8) is correct only for large r as is seen from the assumption for its derivation (compare (9) and $\phi\ll 1$). When ϕ becomes comparable with or larger than unity, Eq. (8) does not hold. If such a region of small r does not play any role for the behavior of the system, the expression (8) can be used as a non-relativistic approximation of W given by (5) for all values of r because of involving no r^{-3} term. Such a possibility was once imagined by Bethe,[6] but finally discarded by him.

From this consideration we find that a r^{-3} term and divergent integrals are due to the unsuitable non-relativistic form of W_{BA} in the customary method. In the actual calculation, the sum in (2) includes relativistic states whereas W_{BA} is non-relativistic. Therefore the non-relativistic form of W_{BA} should be so determined that the contribution from a relativistic region in the sum of (2) does not play any role. The method of the present letter satisfies this requirement, whereas the customary method does not. The difficulties of the latter come from the relativistic region in the sum of (2).

[1] C. Møller and L. Rosenfeld, Kgl. Danske Vid. Sels. Math.-Fys. Medd. **17**, No. 8 (1940).
[2] J. Schwinger, Phys. Rev. **61**, 387 (1942).
[3] G. Araki, Phys. Rev. **74**, 985, 986 (1948).
[4] E. C. Nelson, Phys. Rev. **60**, 830 (1941).
[5] F. J. Dyson, Phys. Rev. **73**, 929 (1948).
[6] H. A. Bethe, Phys. Rev. **57**, 260 (1940).

Microwave Spectrum of Methyl Bromide

A. Harry Sharbaugh and John Mattern
General Electric Research Laboratory, Schenectady, New York
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We have measured and analyzed the $J=0\rightarrow 1$ rotational transition of $C^{12}H_3Br^{79}$ and $C^{12}H_3Br^{81}$ and found the molecular and nuclear constants to be in good agreement with those reported by Gordy, Simmons and Smith[1] for the corresponding $J=1\rightarrow 2$ transition. The measurements were made with a 2K33 Klystron whose low frequency range was extended by increasing the diameter of the external resonant cavity peculiar to this tube. The Wilson-Hughes stark-modulation scheme was used for detection of the lines. Frequency measurements were made to an accuracy of ± 0.1 Mc by using harmonics of a crystal-controlled oscillator which was zero beat with the Bureau of Standards Station WWV.

Table I. Observed and calculated absorption frequencies for $J, K=(0, 0)\rightarrow(1, 0)$ transition of methyl bromide. Ground vibrational state.

	Frequency	
$F\rightarrow F\pm 1$	Observed For $C^{12}H_3Br^{79}$	Calculated* $\nu_0=19136.73$
$3/2\rightarrow 1/2$	18992.36 mc/sec.	18992.41
$3/2\rightarrow 3/2$	19252.13	19252.19
$3/2\rightarrow 5/2$	19107.97	19107.87
	For $C^{12}H_3Br^{81}$	$\nu_0=19064.40$
$3/2\rightarrow 1/2$	18943.77	18943.80
$3/2\rightarrow 3/2$	19160.82	19160.88
$3/2\rightarrow 5/2$	19040.32	19040.28

* $eQ(\partial^2V/\partial Z^2)=577.3$ Mc for $C^{12}H_3Br^{79}$
$eQ(\partial^2V/\partial Z^2)=482.4$ Mc for $C^{12}H_3Br^{81}$

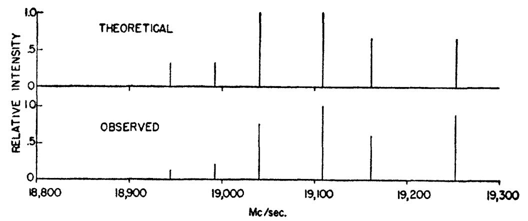

Fig. 1. Microwave absorption spectrum of methyl bromide, $J=0\rightarrow 1$ transition.

Table II. Frequencies of unsplit rotational lines and calculated moments of inertia.

Molecule	Transition	ν_0(mc/sec.)	$I_B\times 10^{40}$ (g-cm^2)	$eQ(\partial^2V/\partial Z^2)$
$C^{12}H_3Br^{79}$	$0\rightarrow 1$	19136.73	87.68_{11}	577.3
$C^{12}H_3Br^{79a}$	$1\rightarrow 2$	38272.40	87.68_{35}	577.0
$C^{12}H_3Br^{81}$	$0\rightarrow 1$	19064.40	88.01_{37}	482.4
$C^{12}H_3Br^{81a}$	$1\rightarrow 2$	38126.97	88.01_{80}	482.0

[a] See reference 1.

In Table I are listed the experimentally observed and calculated absorption frequencies together with the assigned quantum numbers. Second-order quadrupole corrections[2] were of the same order of magnitude as the experimental error and were neglected. Figure 1 shows graphically the comparison between the observed and calculated spectrum.

The frequencies of the unsplit lines (ν_0), moments of inertia (I_B), and quadrupole coupling constants $eQ(\partial^2V/\partial Z^2)$ are compared with those obtained by Gordy *et al.* for the $J=1\rightarrow 2$ transition[1] in Table II. The agreement is seen to be quite satisfactory.

[1] W. Gordy, J. W. Simmons, and A. G. Smith, Phys. Rev. **74**, 243 (1948).
[2] J. Bardeen and C. H. Townes, Phys. Rev. **73**, 627 (1948).

Erratum: Screening and Relativistic Effects on Beta-Spectra

[Phys. Rev. **75**, 264 (1949)]

C. Longmire and H. Brown
Pupin Physics Laboratories, Columbia University, New York, New York

Just prior to publication of this article we noted that in our calculations concerning the screening effect, we had overlooked the energy dependence of the free-electron S wave function at the origin. When this is considered, the radicals in Eqs. (4) and (6) must be inverted, i.e., the correct expression is

$$F_c^{\pm}(Z, W)=F(Z, W\pm D_0)\left(\frac{W-mc^2\pm D_0}{W-mc^2}\right)^{\frac{1}{2}}.$$

Therefore, the numbers in the second column of Table I should be ***divided*** by the numbers in the third column to obtain the fourth column. This modification approximately interchanges the electron and positron corrections. Thus the positron correction is actually the larger of the two. The screening correction for electrons is almost negligible, because of almost complete cancellation of the two factors involved. Even for *RaE* the correction in the Kurie plot amounts to less than (−) one percent at 20 kev. Thus curve (*b*) in Fig. 4 ought to be a horizontal line. The smallness of the effect for electrons was noted correctly long ago by M. E. Rose,[1] with whose result we are now in accord.

We wish to emphasize that correction of this error ***improves*** the over-all agreement between the Cu^{64} spectra obtained by Wu and Albert[2] and the theoretical allowed spectra because of the increased magnitude of the screening correction for positrons, particularly at the higher energies. It is nevertheless true that the relative smallness of the deviations found by Wu and Albert, compared to those previously reported, does not depend to an appreciable extent on whether the screening correction is applied, as may be seen from the uncorrected curves of these authors.

[1] M. E. Rose, Phys. Rev. **49**, 727 (1936).
[2] C. S. Wu and R. D. Albert, Phys. Rev. **75**, 315 (1949). See also C. S. Wu and R. D. Albert, this issue.

Electronic Eigenfunctions in Ionic Crystals

Reikichi Nosawa
Physics Department, Tokyo Institute of Technology, Tokyo, Japan
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The Wigner-Seitz method of approximation, developed further by Slater and others, has proved very useful when applied to the problem of valence-electrons in metals. It must be noted, however, that the approximation thus far used is valid only in the case where the contribution of surrounding cells to the potential can be assumed zero on account of high symmetry and essentially neutral electrical nature, as in metals having cubic structure. The method is, therefore, applicable strictly neither to the ionic crystal where the neutral nature of surrounding cells is lost, nor to the