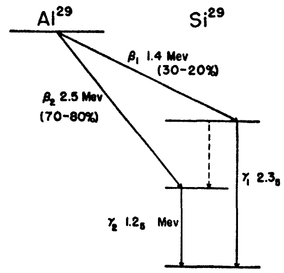

FIG. 2. Decay scheme of Al^{29}.

the sequence of similar nuclei. It seemed interesting, therefore, to redetermine the decay scheme of Al^{29}.

Al^{29} was produced by bombarding Mg with the 18-Mev α-particles of the Purdue cyclotron. The absorption curves for single counts and $\beta\gamma$-coincidences were obtained by following, for each absorber, the decay curve together with that for a standard absorber and analyzing them into the 2.30-min. and 6.56±0.06-min. periods of Al^{28} (from Mg^{25}) and Al^{29}. Contrary to previous measurements,[8] the absorption curve shows a γ-background. The analysis indicates a complex β-spectrum, consisting in the simplest case of two components of 2.5 Mev and 1.4 Mev (∼30 percent). All β-transitions are followed by γ-radiation: the ratio of $\beta\gamma$-coincidences to single β-counts is constant from about 200 mg/cm² to at least 1000 mg/cm²; its value indicates (after calibration of the arrangement with Na^{24} that the 2.5-Mev β-spectrum is followed by a γ-ray of 1.2±0.2 Mev. For thinner absorbers $N_{\beta\gamma}/N_\beta$ increases slightly, corresponding to the higher γ-energy coupled with the low energy partial β-spectrum. In addition, the γ-ray energies were measured directly by coincidence absorption of the secondary electrons. Two components of 2.3_5 Mev (∼20 percent) and 1.2_5 Mev can be distinguished.

The decay scheme compatible with these data is given in Fig. 2. Both β-transitions are allowed with $ft=1.6\cdot 10^5$ (2.5 Mev) and $2\cdot 10^4$. Cascade γ-radiation, with an intensity less than 10 percent, cannot be excluded. While the decay scheme might be subject to modifications, the total disintegration energy of 3.75±0.25 Mev seems well established. The lower curve of Fig. 1 shows the decay energies of the nuclei N^{17} to S^{35}. The value for Al^{29} is seen to be about 1.7 Mev higher than expected from the neighboring members of the sequence. This may be interpreted as confirming evidence for an especially stable structure of Si^{28} (and Si^{29}), due possibly to a shell closure at $Z=14$.

* Supported by the ONR.

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On a New Conception of the Interaction between Charges and Electromagnetic Field

LOUIS DE BROGLIE
Institut Henri Poincaré, Paris, France
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FOR the past fifteen years, I have tried to develop a form of quantum electrodynamics which I have named "wave mechanics of the photon." This theory has been set forth in several books.[1] In 1935, at the beginning of these investigations, a hypothesis was emphasized in order to avoid an infinite value for the self-energy of a charged particle. The wave mechanics of the photon discriminates between the coordinates of the charge X, Y, Z (vector $\mathbf{R}$) and the variables of the field x, y, z (vector $\mathbf{r}$). For the interaction between charge and field in a frame of reference in which the particle is at rest, the following expression is found:

$$\int\int V(\mathbf{r})\rho(R)\delta(\mathbf{r}-\mathbf{R})d\mathbf{r}d\mathbf{R}, \tag{1}$$

where δ is the Dirac function. This expression gives, as usual, an infinite value for the proper energy of the charge. I have suggested[2] that we might replace the δ-function in (1) by a function $f(\mathbf{r}-\mathbf{R})$ in the form of a very sharp, but not infinitely sharp, needle. This substitution, which would express a kind of uncertainty for the point of application of the field on the charge, would lead to a finite value of the self-energy.

Unfortunately, this idea meets difficulties from the viewpoint of relativistic invariance. Let us adopt the multi-temporal theory which introduces distinct times t and T for the field and the charge. The correct relativistic variance of the expression (1) is shown by the fact that one can write it in *any* reference frame:

$$\left|\int\int\sum_{\mu=1}^{4}A_\mu(x,y,z,t)j_\mu(X,Y,Z,T)\left(-\frac{1}{4\pi c}\frac{\partial D(0)}{\partial t}\right)\times dxdydzdXdYdZ\right|_{t=T}, \tag{2}$$

where $D(0)$ is the form for $k_0=0$ of the invariant singular function

$$D(k_0)=-\frac{1}{2\pi^2}\int\frac{\sin[kc(t-T)-\mathbf{k}(\mathbf{r}-\mathbf{R})]}{k}d\mathbf{k} \tag{3}$$

with $k^2=|\mathbf{k}|^2+k_0^2$. Indeed, for all values of k_0

$$\left|\frac{1}{c}\frac{\partial D(k_0)}{\partial t}\right|_{t=T}=-4\pi\delta(\mathbf{r}-\mathbf{R}). \tag{4}$$

But if we will substitute for the δ-function in (1) a function in the form of a sharp needle, we cannot see how the correct relativistic variance can be preserved.

I have been recently lead to a new investigation of this problem by recent papers of R. P. Feynman. Guided in my choice by the suggestions of Feynman and others,[3] I will start from the following form of Wentzel potentials (in non-rationalized units):

$$A_\mu(x,y,z,t,X,Y,Z,T)=-\epsilon\int_{-\infty}^{X_\mu}[D(0)-D(k_0)]dX_\mu, \tag{5}$$

ϵ being the electric charge of the particle. Then we have

$$\Box A_\mu=-k_0^2\epsilon\int_{-\infty}^{X_\mu}D(k_0)dX_\mu;\quad \Sigma_\nu\frac{\partial A_\nu}{\partial x_\nu}=\epsilon[D(0)-D(k_0)]. \tag{6}$$

Starting from the usual definition of the field in terms of potentials, we can easily find for the "electric field" $\mathbf{E}$ the relation

$$\mathrm{div}\mathbf{E}=-\frac{\epsilon}{c}\frac{\partial}{\partial t}[D(0)-D(k_0)]+k_0^2\epsilon\int_{-\infty}^{cT}D(k_0)d(cT). \tag{7}$$

If one makes $T=t$ (equalization of the times), one finds by (4)

$$\mathrm{div}\mathbf{E}=k_0^2\epsilon\left|\int_{-\infty}^{cT}D(k_0)d(cT)\right|_{T=t}. \tag{8}$$

The calculation in a frame of reference where the charge is at rest gives

$$\left|\int_{-\infty}^{cT}D(k_0)d(cT)\right|_{T=t}=\frac{\exp[-k_0|\mathbf{r}-\mathbf{R}|]}{|\mathbf{r}-\mathbf{R}|}. \tag{9}$$

Therefore, we see by (8) and (9) that we can replace the point charge ϵ localized at $\mathbf{R}$ by a continuous distribution of total charge ϵ with the density

$$\sigma=\epsilon f(\mathbf{r}-\mathbf{R})=\frac{\epsilon k_0^2}{4\pi}\frac{\exp[-k_0|\mathbf{r}-\mathbf{R}|]}{[\mathbf{r}-\mathbf{R}]}. \tag{10}$$

It is noteworthy that σ has the form of a Yukawa potential. Moreover, one can easily verify that the electrical distribution (10) gives at the point $\mathbf{r}$ ($\mathbf{R}$ being at the origin) the scalar potential

$$V(r)=\epsilon\frac{1-\exp[-k_0r]}{r} \tag{11}$$

equal to the difference of a Coulomb potential and a Yukawa potential. We are then lead to substitute in (1) the function $f(\mathbf{r}-\mathbf{R})$ defined by (10), and this is consistent with our 1935 assumption.

We can summarize our conclusions by saying that the interaction between electrical charges takes place with the simultaneous intervention of a photon field and a meson field, each charge reacting in an opposite manner with these two fields (more precisely, the meson charge η is equal and opposite to the electric charge ϵ). If we abstract the meson field, we are obliged to replace the point charge by the continuous distribution (10).

We have obtained our results by the multi-temporal theory. Now, using the equations of the wave mechanics of the photon, we will find them again in a very instructive manner.

In the general theory of particles of spin 1 (mere generalization of the wave mechanics of the photon) one finds[4] that the "electric field" of the particle with a proper mass $m_0=hk_0/2\pi c$ verifies the equation

$$\mathrm{div}\mathbf{E}k_0=-k_0^2Vk_0, \tag{14}$$

Vk_0 being the corresponding "scalar potential." The wave mechanics of the photon supposes that the photon possesses an extremely feeble, but not quite zero proper mass $\mu_0=h\gamma/2\pi c$, so that we have

$$\mathrm{div}\mathbf{E}_\gamma=-\gamma^2V\gamma. \tag{15}$$

Let us imagine a field which is the superposition of a photon field with a constant γ and a meson field with a constant $k_0\gg\gamma$. We will find for the total "electric field"

$$\mathrm{div}\mathbf{E}=\mathrm{div}\mathbf{E}_\gamma+\mathrm{div}\mathbf{E}k_0=-\gamma^2V_\gamma-k_0^2Vk_0. \tag{16}$$

The wave mechanics of the photon proves that the photon potential created around it by a charge placed at the origin of coordinates is[5]

$$V_\gamma=\epsilon\frac{e^{-\gamma r}}{r}-\frac{\epsilon}{\gamma^2}\delta(r). \tag{17}$$

The first term is a quasi-Coulomb potential which for $\gamma\rightarrow 0$ gives the usual Coulomb potential. I have named the second term, which is characteristic of the wave mechanics of the photon, the "coincidence potential."

For the meson potential, we have similarly

$$Vk_0=\eta\frac{e^{-k_0r}}{r}-\frac{\eta}{k_0^2}\delta(r) \tag{18}$$

and we assume as before that $\eta=-\epsilon$. Then, substituting (17) and (18) in (16), we find

$$\mathrm{div}\mathbf{E}=-\epsilon^2\frac{e^{-\gamma r}}{r}+\epsilon\delta(r)+k_0^2\epsilon\frac{e^{-k_0r}}{r}-\epsilon\delta(r)$$
$$=4\pi\epsilon\left[\frac{k_0^2}{4\pi}\frac{e^{-k_0r}}{r}-\frac{\gamma^2}{4\pi}\frac{e^{-\gamma r}}{r}\right]. \tag{19}$$

We can therefore overlook the meson field if we replace the point charge by a continuous distribution

$$\sigma(r)=\frac{\epsilon}{4\pi}\left[k_0^2\frac{e^{-k_0r}}{r}-\gamma^2\frac{e^{-\gamma r}}{r}\right]. \tag{20}$$

For $\gamma\rightarrow 0$, we find again the expression (10).

We must point out how the terms $\delta(r)$ disappear in (19). It is the presence of the coincidence potentials in (17) and (18) which makes possible the elimination of these terms. Now, in its usual form, the quantum electrodynamics assumes that the proper mass of the photon is rigorously equal to zero ($\gamma=0$) and consequently ignores the coincidence potential in (17). The foregoing argument shows that, of course, the coincidence potential plays a very important part, and this seems to supply a serious reason for believing that the proper mass of the photon is not quite rigorously zero.

Finally, let us note that a many-meson field could be considered in the interaction, but only if the sum of all the meson charges η_i is equal to $-\epsilon$. This relation between the charges of a particular particle could be of importance in the theory of nuclear fields.[6]

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[3] R. P. Feynman, Phys. Rev. **74**, 939 (1948). A complete bibliography of the question is found in this paper.
[4] L. de Broglie, *Mécanique ondulatoire du photon*, see reference 1, p. 38.
[5] See reference 4, p. 172.
[6] If the relation is valid for a neutral particle, the sum of all its meson charges must be zero.

Atmospheric Ozone Temperatures As Determined by Spectroscopic Observations from Below

William M. Benesch
Physics Department, Johns Hopkins University, Baltimore, Maryland
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The moon at dichotomy provides a light source with a continuously variable surface temperature suitable for measuring the effective temperature of atmospheric ozone in the manner of the line reversal method for flame temperatures. If a relatively cold area of the lunar light source is focused on a spectrometer slit, the 9.6μ ozone band will appear as an emission band when that region of the spectrum is scanned. On the other hand, if a relatively hot strip of the moon is used, the ozone band will appear as an absorption band. That lunar strip which fails to reveal the ozone band has the same temperature, radiatively speaking, as the ozone. I am indebted to Professor Strong for suggesting the use of this method.

To facilitate observations, a hot area on the moon was used at all times and adjusted to cover just that fraction of the slit which would "wash out" the ozone band. This modification makes observations possible using the full moon or even the sun. The ozone temperatures obtained were determined from spectral intensities on the long wave-length side of the 9.6μ band when the slit was so covered that the band was "washed out." These intensities were converted to equivalent temperature by substituting local calibrating black bodies for the moon in front of the moon-follower mirror.

Figure 1 shows three typical spectra. They were taken with (a) excess, (b) matching, and (c) deficient lunar energy.

The following ozone temperatures were obtained in this manner:

April 9, 1949	−47°C
April 11, 1949	−46°C
May 8, 1949	−43°C.

It is to be noted that the precision to which temperatures could be read in these observations was considerably less than is indicated by the consistency of the data above, owing to the inferior sensitivity and resolving power of the dispersing and detecting system. Furthermore, these "effective ozone temperatures," however carefully measured, are restricted in meaning by the effect of pressure broadening on the individual lines of the band. That is, the results obtained here are only representative of the ozone as seen from below, from which aspect the lower layers tend to obscure those at greater altitudes.

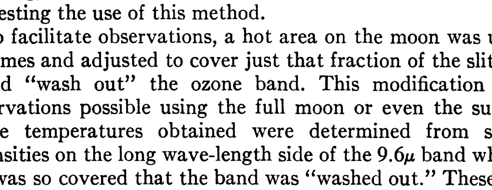

Fig. 1. Three typical spectra with (a) excess, (b) matching, and (c) deficient lunar energy.