International Journal of Theoretical Physics, Vol. 26, No. 10, 1987

Spin-Dynamical Theory of the Wave-Corpuscular Duality

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Received May 19, 1987

The assumption that translations of the electron are accompanied by spin precession enables a deterministic description of electron diffraction and quantization of atomic systems. It is shown that the electromagnetic field of the precessing electron is responsible for modulation of the beam intensity of an electron scattered from a system of charges and for mechanical stability of the orbital motion of electrons in bound states.

1. INTRODUCTION

Since the Uhlenbeck–Goudsmit discovery of spin in 1925, it has been assumed, following a rough macroscopic analogy, that for free translations of the electron its spin axis remains firmly oriented in space. A direct consequence of this was the conclusion that gyromagnetic properties of the electron play a negligible role in shaping electron orbits in atomic systems (Sommerfeld, 1951). Although such a conclusion was methodologically inconsistent (it contradicts the well-known fact that atomic energy levels and electron diffraction phenomena depend explicitly upon the Planck constant h and, therefore, may be directly related with spin properties of the electron), this point of view was commonly accepted. As a result, the development of a deterministic theory of the atom was almost completely stopped.

Positive results during the last two decades in the application of classical physics to the description of atomic collision problems (Gryziński 1965, 1975; Burges and Percival, 1968; Vriens, 1970; Bates, 1978; Grujić *et al.*, 1983, Gryziński and Kunc, 1986) have shown that the concept of a localized electron moving along a well-specified orbit should be considered as a physical reality.

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0020-7748/87/1000-0967$05.00/0

This reality had, however, two great enigmas: quantization of atomic orbits and electron diffraction. Much work has been done on these puzzles (e.g., Bohm, 1952; de Broglie, 1953; Bunge, 1967; Janossy, 1952; Lande, 1960). Not long ago, it was shown that the Kepler problem in the presence of velocity-dependent periodic perturbations has a discrete spectrum of stable orbits (Gryziński, 1980). This spectrum was found to be identical with the spectrum of atomic energy levels providing the frequency of periodic perturbations is proportional to the kinetic energy of the electron and one-half the Planck constant is the coefficient of proportionality. From this fact the hypothesis was developed that translations of the electron are accompanied by precession of the spin axis.

2. TRANSLATIONAL PRECESSION OF THE SPIN

Ascribing gyromagnetic properties to the electron brings the need of formulating the laws of motion for spin coordinates, that is, for the two angles defining the orientation of the electron spin axis in space. The classical theory of fields predicts the existence of electron spin axis precession in the presence of electromagnetic fields (e.g., Barut, 1964), but it leaves completely open the question of transient effects in the behavior of spin, the theory of which should be formulated in second time derivatives of spin coordinates, and says nothing about the behavior of spin in free translations of the electron. In view of the above and given various paradoxes in the behavior of the electron, such as the Aharanov-Bohm (Aharanov and Bohm, 1959), the problem seems to be open for discussion.

Let us assume, therefore, having in view, for instance, some hydrodynamic analogy, and in accordance with the conclusion of Gryziński (1980), that translations of the electron are accompanied by precession of the spin axis around the velocity vector $\mathbf{v}$ and that the angular velocity of the precession is equal to the kinetic energy of the electron divided by its intrinsic angular momentum. Thus, the postulated law of motion for the spin of the free electron is

$$\frac{d\mathbf{s}}{dt} = \left(\frac{mv}{\hbar}\right)[\mathbf{s} \times \mathbf{v}] \tag{P.1}$$

where $\mathbf{s}$ is the unit vector directed along the electron spin axis. If one denotes the angle defining the orientation of the spin axis in the plane perpendicular to the velocity vector $\mathbf{v}$ by ψ_s, then on the grounds of (P.1) one has

$$\frac{d\psi_s}{dt} = \frac{mv^2}{\hbar} \tag{2}$$

or alternatively

$$\frac{d\psi_s}{dl} = \frac{mv}{\hbar} \tag{3}$$

where dl is the element of the electron path.

Let us denote now by λ the distance in which the cyclic coordinate ψ_s changes the value by 2π. Then the orientation of the spin axis after traversing the distance λ is the same as at the beginning of the path. Let us integrate equation (3) along this path. As a result, we obtain

$$\int_0^{2\pi} d\psi_s = \left(\frac{m}{\hbar}\right) \int_0^{\lambda} v\, dl \tag{4}$$

If the electron moves with a constant velocity, the spin wavelength equation (4) assumes the form of the famous de Broglie wave postulate:

$$\lambda = \frac{h}{mv} \qquad \text{(de Broglie postulate)} \tag{5}$$

This result shows the equivalence of the postulate (P.1) and the de Broglie postulate. However, now the wavelength has a well-specified meaning: it is the distance traveled by the electron during one revolution of the spin axis. Moreover, now the "wave" relation, when formulated in the differential form, opens new possibilities. In view of the experimental fact that the electron has magnetic moment $\boldsymbol{\mu}$,

$$\boldsymbol{\mu} = e\frac{\hbar}{mc}\mathbf{s} \tag{6}$$

which produces the field

$$\mathbf{H} = \nabla \times \mathbf{A}_\mu \qquad \text{where} \quad \mathbf{A}_\mu = \frac{\boldsymbol{\mu} \times \mathbf{r}}{r^3} \tag{7}$$

the postulate (P.1) has a great significance for the whole dynamics of the moving electron.

Now the dipole magnetic field of the moving electron is a periodically varying function of time, and, as follows from the basic relation of the Maxwell theory

$$\nabla \times \mathbf{E} = -\frac{1}{c}\frac{\delta \mathbf{H}}{\delta t} \tag{8}$$

the moving electron has a periodic varying spin component of the electric field:

$$\mathbf{E}_s^{\text{spin}} = -\frac{1}{c}\left[\frac{\dot{\boldsymbol{\mu}} \times \mathbf{r}}{r^3}\right] \tag{9}$$

The field of the moving electron in the presence of spin translational precession is therefore

$$\mathbf{E}_{q,s}=\frac{-e}{r^2}\left\{\left(\hat{r}+\left(\frac{v}{c}\right)^2[(\mathbf{s}\times\hat{\mathbf{v}})\times\hat{\mathbf{r}}]\right\}, \tag{10}$$

$$\mathbf{H}_s=-e\lambda\!\!\!^{-}_C[\nabla\times(\mathbf{s}\times\mathbf{r}/r^3)] \tag{11}$$

where $\lambda\!\!\!^{-}_C$ is the electron Compton wavelength, and $\mathbf{s}(t)$, given by equation (P.1), is a periodic function of time. For a free electron $\mathbf{s}(t)$ can be written in the following way:

$$\mathbf{s}(t)=(\mathbf{s}\cdot\hat{\mathbf{v}})\hat{\mathbf{v}}+[(\mathbf{s}\times\hat{\mathbf{v}})\times\hat{\mathbf{v}}]\exp(i\omega_s t) \tag{12}$$

which clearly shows the periodic character of the spin field.

In view of the above it is reasonable to expect, during the motion of an electron in a system of periodically situated charges or current elements, various resonances introducing periodic modulations in the intensity of scattered electrons characteristic for diffraction and interference.

3. DIFFRACTIONAL SCATTERING

To explain the essence of the electron diffraction (interference) phenomena, let us consider scattering of a single fast electron on identically charged collinear lines (see Fig. 1). Under the assumption that the potential produced by the scattering system is appreciably smaller than the potential accelerating the electrons, the problem can be effectively solved on the basis of the co-called small-angle approximation, which is a satisfactorily precise method of analysis for most electron diffraction experiments. In the approximation considered the scattering angle is given by

$$\tan\vartheta\simeq\frac{1}{mv_0}\int_{-\infty}^{+\infty}\hat{\mathbf{k}}\cdot\mathbf{F}(x=\text{const},\,y=\text{const},\,z=v_0t)\,\mathrm{d}t \tag{13}$$

where $\hat{\mathbf{k}}$ is the unit vector perpendicular to the initial velocity of the electron and $\mathbf{F}$ is the interaction force of the electron with the scattering system. If the scattering system is formed by a set of collinear charged lines parallel to the x axis and crossing the yz plane at points y_i, z_i and carrying per unit of length the charge $(q_l)_i$, then

$$\mathbf{F}=\sum_i(q_l)_i\int_{-\infty}^{+\infty}\mathbf{E}_{q,s}(x;\,y_i,\,z_i)\,dx \tag{14}$$

Denoting the distance between the moving electron and the ith charged line by ρ_i,

$$\rho_i^2=(y_i-y_0)^2+(z_i-z_0)^2 \tag{15}$$

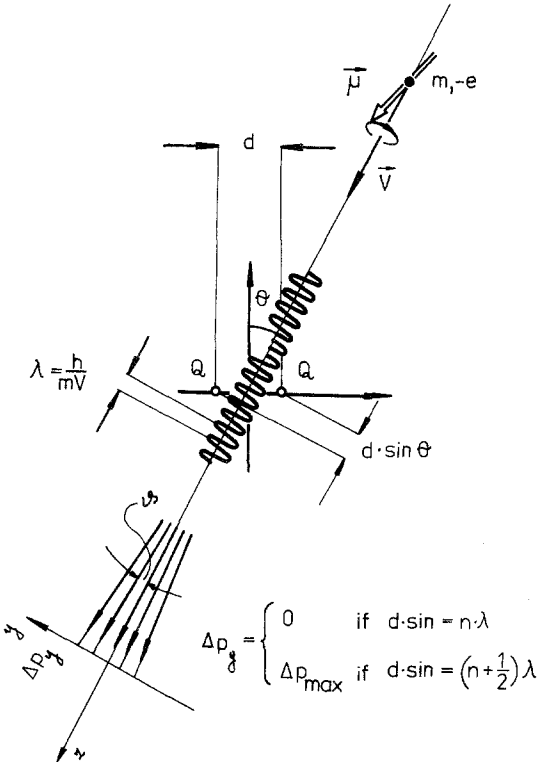

Fig. 1. Whether or not fast electrons are scattered by the two charged centers depends upon the electron spin wavelength and the "effective" distance between centers, which, for electrons moving along a line inclined to the dipole axis at an angle θ, is equal to $d \sin \theta$. This fact forms the essence of a large number of diffraction phenomena.

and taking into account equation (10), we find

$$\mathbf{F} = -e \sum_i \frac{(q_l)_i}{\rho_i} 2\left\{\hat{\rho}_i + \left(\frac{v_0}{c}\right)^2 [(\hat{\mathbf{s}} \times \hat{\mathbf{v}}_0) \times \hat{\boldsymbol{\rho}}_i]\right\} \tag{16}$$

Therefore,

$$\tan \vartheta = -\frac{2}{mv_0^2} e \sum_i (q_l)_i \int_{-\infty}^{+\infty} \frac{1}{\rho_i} \left[\hat{\boldsymbol{\rho}}_i \cdot \hat{\mathbf{k}}_y - \left(\frac{v_0}{c}\right)^2 (\hat{\boldsymbol{\rho}}_i \cdot \hat{\mathbf{k}}_z)(\hat{\mathbf{s}} \cdot \hat{\mathbf{k}}_y)\right] dz \tag{17}$$

According to equation (12), the projection of the spin vector $\hat{\mathbf{s}}$ on the direction $\hat{\mathbf{k}}_y$ in the approximation considered is given by

$$\hat{\mathbf{s}} \cdot \hat{\mathbf{k}}_y = s_y^0 \cos(z/\lambda + \psi_0) \tag{18}$$

Performing an integration over z, we finally obtain

$$\tan\vartheta = -\frac{2\pi}{mv_0^2} e \sum (q_l)_i \left[\operatorname{sgn}(y_i - y_0) - \left(\frac{v_0}{c}\right)^2 s_y^0 \sin\left(\frac{z_i}{\lambda} + \psi_0\right)\right] e^{-|y_0 - y_i|/\lambda} \quad (19)$$

Now, the intensity of electrons scattered at the angle ϑ can be quite easily calculated. Using the δ-function formalism we have

$$I(\vartheta, \psi_0) = \int_{-\infty}^{+\infty} I(y_0)\, \delta[\vartheta - \vartheta(y_0, \psi_0)]\, dy_0 \quad (20)$$

where $\vartheta(y_0, \psi_0)$ represents the solution of equation (19) and $I(y_0)$ is the intensity of the incident beam. Since the effective intensity is the average of scattering events with various values of the phase angle ψ_0, we have

$$I(\vartheta) = \frac{1}{2\pi} \int_0^{2\pi} I(\vartheta, \psi_0)\, d\psi_0 \quad (21)$$

To carry out the calculations to completion one must specify precisely the scattering system. Prior to the analysis of the particular case, however, it is worth noting that in the linear scattering theory with some symmetry in the distribution of the charge in the scattering system the Coulomb interaction terms cancel each other and the scattering is entirely determined by the spin field. In particular, as follows from equation (19), the Coulomb term vanishes if the total charge situated to the left of the coordinate y_0 and the total charge situated to the right of the coordinates y_0 are equal (in particular both may be zero).

Let us assume now that the scattering system is formed by two identically charged lines separated by a distance d and the incident beam is inclined to the normal of the plane containing the lines by the angle θ; then

$$y_i = (-1)^i \frac{d}{2} \cos\theta, \qquad z_i = (-1)^i \frac{d}{2} \sin\theta, \qquad i = 1, 2 \quad (22)$$

and the scattering formula assumes the form

$$\tan\vartheta = -4\pi \left(\frac{eq_l}{mc^2}\right) s_y^0\, e^{-(d\cos\theta)/\lambda} \left\{ \left[\cos\left(\frac{d\sin\theta}{\lambda}\right) \sin\psi_0 \operatorname{ch}\left(\frac{y_0}{\lambda}\right) + \sin\left(\frac{d\sin\theta}{\lambda}\right) \cos\psi_0 \operatorname{sh}\left(\frac{y_0}{\lambda}\right)\right]\right. \quad (23)$$

Introducing the above into equation (20), one can calculate the intensity of the scattered beam. Assuming that the intensity of the incident beam is

uniform across the beam, we have

$$I(\vartheta) = (I_0\lambda)\frac{1}{2\pi}\int_0^{2\pi}\int_{-(d\cos\theta)/\lambda}^{+(d\cos\theta)/\lambda} \delta[\tan\vartheta - \tan\vartheta_0(\cos D\sin\psi_0 \operatorname{ch}\xi + \sin D\cos\psi_0 \operatorname{sh}\xi)]\, d\psi_0\, d\xi \tag{24}$$

where

$$\tan\vartheta_0 = 4\pi\left(\frac{eq_l}{mc^2}\right)s_y^0\, e^{-(d\cos\theta)/\lambda} \tag{25}$$

and

$$D = \frac{d\sin\theta}{\lambda} \tag{26}$$

Performing an integration over ψ_0, we obtain

$$I(\vartheta) = (I_0\lambda)\frac{2}{\pi}\int_{-(d\cos\theta)/\lambda}^{+(d\cos\theta)/\lambda} \frac{d\xi}{\sqrt{[(\operatorname{ch}^2\xi - \sin^2 D)\tan^2\vartheta_0 - \tan^2\vartheta]^{1/2}}} \tag{27}$$

With the assumption $d(\cos\theta)/\lambda \gg 1$, the integration can be effectively carried out; the result is

$$I(\vartheta) \simeq \left(\frac{I_0\lambda}{\tan\vartheta_0}\right)K\left(\left[\sin^2 D + \left(\frac{\tan\vartheta}{\tan\vartheta_0}\right)^2\right]^{1/2}\right) \tag{28}$$

where K is the complete integral of the first kind.

It follows from the above that the width and intensity of the scattering line are respectively given by

$$\tan\vartheta_{\max} = \tan\vartheta_0\cos\frac{d\sin\theta}{\lambda} \to 0 \qquad \left(\text{if} \quad \cos\frac{d\sin\theta}{\lambda} \to 1\right) \tag{29}$$

and

$$I(\vartheta = 0) = \frac{I_0\lambda}{\tan\vartheta_0}K\left(\sin\frac{d\sin\theta}{\lambda}\right) \to \infty \qquad \left(\text{if} \quad \sin\frac{d\sin\theta}{\lambda} \to 1\right) \tag{30}$$

The structural parameter d/λ is therefore the factor that determines the transmission of the system. Electrons can pass through the system unscattered if

$$\sin\theta_n = n\lambda/d \tag{31}$$

It follows from the above that electrons penetrating a system of periodically situated charges can propagate freely (not being scattered) in the directions determined by equation (31). The directions of electron propagation correspond to individual spots on the diffraction picture produced by electrons penetrating the crystals.

In the case of electrons moving through a slit or through a system of slits (as in the case of Young type experiments) the situation is almost identical. To realize this one must take into account that the distribution of the charge in a conducting plane with a slit is highly peaked at the edges of the slit. According to the classical theory of the field this distribution is given by

$$\sigma \sim x/[x^2-(d/2)^2]^{1/2} \qquad \text{if} \quad x > d/2 \tag{32}$$

where d is the width of the slit and x is the distance to the center of the slit. Since $\sigma \to \infty$ at $x \to d/2$, therefore the field in the slit may in first approximation be identified with the field of two charged lines placed at the edges of the slit and the whole scattering problem may be reduced to the case discussed above.

In view of relation (29) and (30), it is evident that the shadow of a slit illuminated with electrons emitted from a point source is not uniform, since

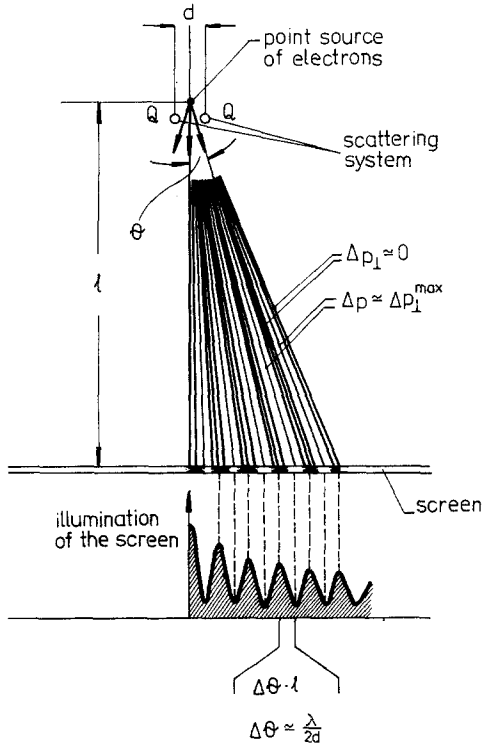

Fig. 2. The diffraction of electrons. Depending upon the value of the angle θ, electrons are scattered or not. As a result, illumination of the screen becomes modulated.

at angles θ_n given by equation (31) electrons proceed without deviation, and at angles $\theta_{n\cong 1/2}$ given by

$$\sin \theta_{n+1/2} = (n+\tfrac{1}{2})\frac{\lambda}{d} \tag{33}$$

transmission of the electrons as a result of scattering is strongly lowered. As a result a diffraction picture as shown in Fig. 2 is formed.

The diffraction of electrons at the edge of an object has the same origin; however, in the case of a nonconducting object, one must keep in view the fact that matter consists of positively and negatively charged particles and that the charge neutrality locally (i.e., at the boundary of the objects) is not preserved.

4. DYNAMICS OF QUANTIZATION

In quantum wave mechanics the quantized states are determined by sets of rules (the Schrödinger equation may be considered as a rule formulated in differential form) and the dynamics of transitions between the quantized states is not considered at all.

In our case, identification of the "wave" field of the electron forms the ground for the development of a dynamical theory of quantization with a precise description of the transition process between quantized states. Here we present the general aspects of the problem.

Formally the problem we have to solve first is the problem of the mechanical stability of electron motion in the presence of an electromagnetic spin field.

In atomic systems the electromagnetic spin field is in general smaller than the Coulomb field by $(Z\alpha)^2$. Nevertheless, it cannot be neglected, since in bound states these small periodic perturbations, after a long enough time, through accumulations of small amounts of energy and momentum, may change appreciably the parameters of the orbit.

The stability of mechanical systems in the presence of perturbations is one of the most important problems of classical mechanics. It was shown some time ago that for some class of perturbations the mechanical system has a discrete spectrum of stable solutions. According to Czetaev (1962), for the class of perturbations satisfying the set of equations

$$\psi = A \exp(ikS) \tag{34}$$

$$\sum_{i,j} \frac{\delta}{\delta q_i}\left(g_{ij}\frac{\delta A}{\delta q_j}\right) = k^2 W \tag{35}$$

$$\sum_{i,j} g_{ij}\frac{\delta A}{\delta q_j}\frac{\delta S}{\delta q_i} = 0 \tag{36}$$

the condition of mechanical stability is given by

$$\sum_{ij} \frac{\delta}{\delta q_i}\left(g_{ij}\frac{\delta\psi}{\delta q_j}\right)+2k^2(\varepsilon-u)=0 \tag{37}$$

where U is a potential energy function defining the unperturbed problem, W is the perturbation potential, S is the complete integral of the Hamilton-Jacobi equation for the unperturbed motion, and $q_{i(j)}$ are generalized coordinates. It is well known that the differential equation of the form given above has stationary solutions for particular values of ε only; in another word, energy corresponding to the stationary motion is quantized.

Let us look now at the mechanical stability equations (34)-(37) from the point of the view of the translational precession law formulated above.

For a particle moving in the field, S is given by

$$S=-\varepsilon t+\int_0^l \mathbf{p}\cdot d\mathbf{l} \tag{38}$$

As follows from equations (2) and (3), S/h represents the phase angle describing the difference in azimuthal orientation of the spin vector $\mathbf{s}$ with respect to the velocity vector $\mathbf{v}$ in the case of the particle moving with a constant velocity (constant energy):

$$\psi_s^0=\frac{1}{\hbar}\varepsilon t \tag{39}$$

and in the case of a particle moving in the field $U(r)$

$$\psi_s=\frac{1}{\hbar}\int_0^l \mathbf{p}\cdot d\mathbf{l} \tag{40}$$

Therefore,

$$k=1/\hbar \tag{41}$$

$$\psi\sim\exp i(\psi_s-\psi_s^0) \tag{42}$$

and

$$\Delta\psi+(2m/\hbar^2)(\varepsilon-U)=0 \tag{43}$$

It follows from the above that the Schrödinger equation is identical with the classical mechanical stability equation determining the stability of the electron orbit in the presence of spin perturbations, where the function ψ represents the translational precession of electron spin (real and imaginary parts of ψ represent projections of the spin vector on the two perpendicular axes of the fixed system of coordinates).

If the electron moves in the field of the nucleus, then

$$U=-Ze^2/r, \tag{44}$$

and the values of ε at which the electron motion is stable are the same as given by Bohr's quantum postulate:

$$\varepsilon = -2\pi m Z^2 e^4 / n^2 \hbar^2 \tag{45}$$

If the perturbation force is known, one can derive the spectrum of stable electron orbits explicitly by direct solution of the perturbed Kepler problem. According to the general method of perturbation calculus, the Kepler orbit may be stable if

$$\delta\varepsilon = \int_0^T (\mathbf{F}_p \cdot \mathbf{v})\, dt = 0 \tag{46}$$

where $\mathbf{F}_p$ is the perturbation force and $\mathbf{v}$ is the orbital velocity, and the integeral is taken along the unperturbed ellipse. Taking into account equation (9), which defines the perturbation force $\mathbf{F}_p$, one can show, performing partial integration, that

$$\delta\varepsilon = -\frac{Ze}{c}(\mathbf{A} \cdot \mathbf{v})\Big|_0^T \tag{47}$$

In the presence of translational precession $\delta\varepsilon$ is zero if the spin axis performs a complete number of revolutions along the whole orbit. This requirement, according to equation (2), is satisfied if

$$\int_0^{n2\pi} d\psi_s = \int_0^T \left(\frac{mv^2}{\hbar}\right) dt \tag{48}$$

The above relation, which is in fact identical with Bohr's postulate of quantization, determines the spectrum of stable Kepler orbits. The spectrum is of course the same as given by equation (45).

The above analysis shows that the postulate (P.1) is the equivalent not only of the de Broglie postulate, but of the Schrödinger postulate as well. As in the case of the de Broglie wave postulate, the translational precession postulate (P.1) discloses the physical essence of the equation postulated by Schrödinger, and by identification of the forces responsible for quantization opens the way to a formulation of the deterministic theory of the atom and a dynamical theory of atomic transitions.

It is worth noting that the dynamical approach to the quantization of electron orbits in the Coulomb field of the nucleus allows one to draw some important conclusions.

The direct way of solving the perturbation Kepler problem shows the peculiar role of the electron orbit with angular momentum equal to zero. In this case, as follows from equation (12), the perturbation force is zero and the motion is particularly stable (the absolute stability is determined by higher order perturbation terms). This is consistent with the conclusion

derived from atomic collision experiments, which led some time ago to the concept of the free-fall atomic model, according to which the ground-state electron orbit is the zero angular momentum orbit (Gryziński 1972, 1973).

Another conclusion that can be drawn on the grounds of postulate (P.1) concerns the dynamics of the transitions of the electron between "quantized" states.

If the binding energy of the electron is different from that given by equation (45), that is, when the condition of periodicity for the spin coordinate is not satisfied, then the electron is accelerated by the rotary electric field. The change of energy at each trip of the electron around the nucleus follows from equation (47) roughly as

$$|\delta\varepsilon| \simeq Ze^2\lambda\!\!\!^{-2}_{C}(L/h)(1/r^3) \tag{49}$$

where L is the angular momentum of the electron in orbital motion. Taking into account that

$$\overline{\left(\frac{1}{r^3}\right)} = \left(\frac{Z}{n}\right)^3 \frac{1}{a_0^3}\left(\frac{h}{L}\right)^3 \tag{50}$$

one can estimate the evolution time of the orbit. The time necessary for the transition of an electron from an orbit of energy E_{n+1} to an orbit of energy E_n is roughly

$$\Delta t_{(n+1)\to n} \simeq \frac{T_0}{(Z\alpha)^2}\left(\frac{L}{h}\right)n^3 \tag{51}$$

where T_0 is the period of motion on the Bohr orbit.

This estimation, which is in rough agreement with measurements, shows that translational precession muct be taken into account if one is to formulate a dynamical theory of atomic transitions.

Investigating the stability of the Kepler orbit in the presence of perturbations, one must keep in view that not only energy, but angular momentum as well must remain constant in time. This means that the following requirement must be satisfied:

$$\delta L = \int_0^T (\mathbf{F}_p \times \mathbf{r}) \cdot \hat{\mathbf{L}}\, dt = 0 \tag{52}$$

The above requirement imposes additional restrictions on the orbital motion of the electrons, which result in "quantization" of the angular momentum.

Inspecting equation (10), (46), and (52), one finds that during the evolution of the electron orbit the changes of energy are always accompanied by changes of the angular momentum (δL is different from zero if only $\delta\varepsilon$ is different from zero). This result explains the well-known spectroscopic fact that Δl in radiative transitions cannot be equal to zero.

In view of all the above, we arrive at the conclusion that radiative transitions in atomic systems are directly connected with the electron spin axis precession, and acceleration is to be regarded as a necessary but not sufficient condition for the existence of radiation.

5. FINAL REMARKS AND CONCLUSIONS

In conclusion, we can state that the wave-corpuscular enigma is the result of an incorrectly formulated law of motion for the electron spin coordinates (for free motion of the electron its spin axis was assumed to be firmly oriented in space). The postulated translational precession law discloses the nature of the wave-corpuscular duality and forms the basis for a causal deterministic description of atomic processes.

The discovered translational precession law specifies the force responsible for the formation of stable orbits and for diffraction phenomena. It allows one to describe the dynamics of transitions between stable orbits—between different "quantized" states—and to formulate the dynamical theory of diffraction.

Against the common conviction that only wave objects can generate a periodically modulated intensity, it was proved that a periodically modulated intensity of scattered electrons can be derived from deterministic laws of motion for the spinning particle, and that the diffraction theory can be formulated in a causal way. The geometrical relations that form the essence of the existing wave theories have a well-specified origin in the interaction of scattered particle with the inhomogeneous distribution of the charged components of the matter (electrons and nuclei) at the boundaries of material bodies.

The fact that quantization appears as a result of perturbations in the Kepler motion explains the success of classical atomic collision theories, which were developed from the assumption that the atom is a well-specified mechanical system with precisely determined electron trajectories.

There are reasons to think that a similar situation exists in the case of nuclear matter, and that the formation of stable configurations is determined by translational precession of nucleons. However, now magnetic spin interactions seem to play the decisive role.

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