

# RELATIVITY AND NONLOCALITY

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The basic physical structure of the relativistic theory of gravitation is discussed. The significant role that the Hypothesis of Locality plays in relativity theory is elucidated via the phenomenon of spin-rotation coupling. The limitations of this hypothesis are critically examined. A nonlocal theory of accelerated observers is presented and some of its observational consequences are described.

## 1 Introduction

The general theory of relativity is in good agreement with experimental data available at present; however, a microphysical understanding of gravity is still missing. There are many different attempts to reconcile general relativity with the quantum theory. The approach adopted here seeks to examine the physical foundations of general relativity in the light of the principles of quantum physics.

To arrive at the basic physical postulates of general relativity, we adopt a measurement-theoretic approach. Therefore, we start with a global inertial frame of reference in Minkowski spacetime with coordinates  $x^\mu = (ct, \mathbf{x})$  and define the standard observers to be the class of inertial observers at rest in this system. Each standard observer carries an orthonormal tetrad  $\lambda^\mu_{(\alpha)}$ , where  $\lambda^\mu_{(0)} = dx^\mu/d\tau$  is the vector tangent to the worldline of the observer and  $\lambda^\mu_{(i)}$ ,  $i = 1, 2, 3$ , are the spatial axes of the observer. Hence  $\lambda^\mu_{(\alpha)} = \delta^\mu_\alpha$  in this case. The electromagnetic field as measured by the standard observers is the projection of the Faraday tensor  $F_{\mu\nu}$  on the tetrad of the observers,

$$F_{(\alpha)(\beta)} = F_{\mu\nu} \lambda^\mu_{(\alpha)} \lambda^\nu_{(\beta)} . \quad (1)$$

Thus the electric and magnetic fields,  $F_{(\alpha)(\beta)} \rightarrow (\mathbf{E}, \mathbf{B})$ , that appear in the standard form of Maxwell's equations are the fields as measured by the standard inertial observers. Imagine now an inertial observer moving uniformly with speed  $v$  along the  $z$ -axis. To find the fields measured by this observer, one can either apply the Lorentz transformation to the Faraday tensor, i.e. transform to the rest frame of the observer, or to the tetrad frame and then project the field on the tetrad frame of the moving observer in accordance with

(1). We adopt the second approach for the sake of convenience and note that for this observer

$$\begin{aligned}\lambda^\mu_{(0)} &= \gamma(1, 0, 0, \beta) \quad , \quad \lambda^\mu_{(1)} = (0, 1, 0, 0), \\ \lambda^\mu_{(2)} &= (0, 0, 1, 0) \quad , \quad \lambda^\mu_{(3)} = \gamma(\beta, 0, 0, 1) \quad ,\end{aligned}\tag{2}$$

where  $\gamma$  is the Lorentz factor and  $\beta = v/c$ . In this way Lorentz invariance (or Poincaré invariance) provides a complete description of phenomena according to inertial observers, since the fundamental laws of physics have all been formulated with respect to such ideal observers.

To describe the measurements of arbitrary observers in Minkowski spacetime, we need to remove two limitations that exist in our treatment thus far. Let us first demand that inertial observers be able to use arbitrary coordinates in Minkowski spacetime just as one uses curvilinear coordinates in Euclidean space. A detailed examination shows that this limitation is purely mathematical and using the elegant tools of tensor calculus one can simply extend the Lorentz-invariant theory of inertial observers to arbitrary admissible coordinates in Minkowski spacetime. We must next remove the restriction that thus far only inertial observers make measurements; that is, we need a prescription for the results of measurements by arbitrary accelerated observers in Minkowski spacetime. To this end, a basic hypothesis is needed to connect the measurements of accelerated observers to those of inertial observers. It is natural to expect such a connection, since the whole observational basis of Lorentz invariance involves measurements made by accelerated observers such as those on the rotating Earth. The assumption that is employed in the theory of relativity is the Hypothesis of Locality, which postulates the local equivalence of an accelerated observer with an instantaneously comoving (hypothetical) inertial observer. The worldline of an accelerated observer is a curved path in Minkowski spacetime; therefore, the hypothesis of locality replaces the curve by the line tangent to the curve at each event. This hypothesis is embodied in the assumption that clocks and rods are locally unaffected by acceleration [1].

The hypothesis of locality is clearly valid in the Newtonian mechanics of point masses, since the accelerated and inertial particles instantaneously share the same *state*  $(\mathbf{x}, \mathbf{v})$ ; therefore, no extra physical assumption is needed in the discussion of inertial effects in Newtonian mechanics. In relativity theory, however, the pointwise character of the hypothesis of locality imposes severe limitations on the notion of distance [2]. Moreover, the hypothesis of locality implies that an accelerated observer in Minkowski spacetime is also endowed with an orthonormal tetrad frame  $\lambda^\mu_{(\alpha)}(\tau)$  that instantaneously coincides with

that of the comoving inertial observer. From

$$\frac{D\lambda^\mu_{(\alpha)}}{d\tau} = \phi_{\alpha}^{\beta} \lambda^\mu_{(\beta)} \quad (3)$$

and the decomposition of  $\phi_{\alpha\beta} = -\phi_{\beta\alpha}$  in terms of translational acceleration  $\mathbf{g}$  and rotational frequency of the spatial frame  $\mathbf{\Omega}$  with respect to nonrotating spatial axes,  $\phi_{\alpha\beta} \rightarrow (\mathbf{g}, \mathbf{\Omega})$ , we arrive at the acceleration lengths  $c^2/g$  and  $c/\Omega$  that characterize the curvature of the worldline [2-4]. Neglecting the instantaneous curvature of the worldline, as demanded by the hypothesis of locality, means that these acceleration lengths are immeasurably large compared to any other relevant scale in the physical phenomenon under consideration. This limitation will be discussed in the following section.

The special theory of relativity is therefore based upon Lorentz invariance and the hypothesis of locality. To explain the measurements of an observer in a gravitational field, Einstein's principle of equivalence is indispensable. According to this heuristic principle, an observer in the gravitational field is locally equivalent to a certain accelerated observer in Minkowski spacetime. In conjunction with the hypothesis of locality, Einstein's principle of equivalence implies that every observer in a gravitational field is locally inertial. The simplest way to connect the local inertial frames of observers in a gravitational field is via the Riemannian curvature of the spacetime manifold. This curvature is then identified with the gravitational field in Einstein's theory, though many extensions and generalizations are possible. Free test particles and null rays thus follow geodesics of the curved spacetime in general relativity. It remains to give the field equations for gravitation; in general relativity, this last step involves the simplest generalization of Newtonian gravitation that is consistent with the spacetime structure. The physical basis of general relativity thus consists of [5]

- (1) Lorentz Invariance,
- (2) Hypothesis of Locality,
- (3) Einstein's Principle of Equivalence,
- (4) Correspondence with Newtonian Gravitation; Field Equations.

The great success of relativistic quantum theory means that we must begin with the hypothesis of locality in our quest for the integration of relativity theory with quantum physics.

## 2 Spin-Rotation Coupling

Classically, motion takes place via classical particles as well as electromagnetic waves. The latter have an intrinsic length scale characterized by their wavelength  $\lambda$ . For most laboratory experiments, however,  $\lambda$  is very small compared

to the usual acceleration lengths, since  $c^2/g \simeq 1$  yr for the acceleration of gravity on the Earth and  $c/\Omega \simeq 28$  AU for the proper rotation frequency of the Earth. To explore the consequences of the hypothesis of locality explicitly, let us consider a simple thought experiment. Imagine an observer following a circle of radius  $r$  about the origin in the  $(x, y)$ -plane of a global inertial frame of reference such that the observer rotates uniformly around the  $z$ -axis with frequency  $\Omega$ . A plane monochromatic electromagnetic wave propagates along the  $z$ -direction. The frequency of the wave according to the standard observers is  $\omega$ . What is the frequency as measured by the rotating observer? Introducing the hypothetical comoving inertial observer as in the hypothesis of locality, we can employ the transverse Doppler effect between the instantaneous inertial frame and the global inertial frame, and conclude that the answer is simply  $\gamma\omega$  due to time dilation. Here  $\gamma = (1 - \beta^2)^{-1/2}$  and  $\beta = r\Omega/c$ . On the other hand, we can apply this hypothesis to the field and use the tetrad frame of the observer,

$$\begin{aligned} \lambda^\mu_{(0)} &= \gamma(1, -\beta\sin\varphi, \beta\cos\varphi, 0), & \lambda^\mu_{(1)} &= (0, \cos\varphi, \sin\varphi, 0), \\ \lambda^\mu_{(2)} &= \gamma(\beta, -\sin\varphi, \cos\varphi, 0), & \lambda^\mu_{(3)} &= (0, 0, 0, 1), \end{aligned} \quad (4)$$

where  $\varphi = \Omega t = \gamma\Omega\tau$ , to find  $F_{(\alpha)(\beta)}(\tau)$  in accordance with (1) and then Fourier analyze this result — a nonlocal operation — to conclude that the frequency measured by the observer is

$$\omega' = \gamma(\omega \mp \Omega). \quad (5)$$

Here the upper sign is for positive helicity radiation (RCP or right circularly polarized) and the lower sign is for negative helicity radiation (LCP or left circularly polarized). Equation (5) has a simple physical interpretation: in RCP (LCP) radiation, the electromagnetic field rotates in the positive (negative) sense along the direction of propagation and hence the observer sees a kind of rotational Doppler effect. Multiplying both sides of equation (5) by  $\hbar$ , we find  $E' = \gamma(E - \boldsymbol{\sigma} \cdot \boldsymbol{\Omega})$ , which indicates that the spin of the photon couples to rotation. This is an example of the general phenomenon of spin-rotation coupling that is supported by experiment [6-8].

Several features of spin-rotation coupling (5) should be mentioned. Writing (5) as  $\omega' = \gamma\omega(1 \mp \Omega/\omega)$ , we note that the result is different from  $\gamma\omega$  that was obtained by the simple application of the hypothesis of locality by a term of the form  $\Omega/\omega = \lambda/(c/\Omega)$ , where  $c/\Omega$  is the acceleration length. Thus  $\omega' = \gamma\omega$  for a null ray with  $\lambda = 0$ ; hence, the Doppler effect is in general strictly valid only in the geometric optics limit [9]. Moreover, if the incident RCP and LCP waves have the same amplitude according to the standard observers,

then their amplitudes measured by the rotating observer will also be the same. Finally, equation (5) is an example of the more general formula  $\omega' = \gamma(\omega - M\Omega)$  for oblique incidence, where  $M$  is a total spin parameter given by  $M = 0, \pm 1, \pm 2, \dots$ , for a scalar or a vector field, while  $M \mp 1/2 = 0, \pm 1, \pm 2, \dots$ , for a Dirac field. Let us note that  $\omega' = 0$  for  $\Omega = \omega/M$ , i.e. an observer can stand completely at rest with respect to the radiation field by a mere rotation. This is particularly clear in (5) for an observer that rotates with frequency  $\Omega = \omega$  with respect to an incident RCP wave. This basic difficulty, which does not arise in the Doppler effect, leads us to re-examine the physical basis of the hypothesis of locality.

### 3 Nonlocality

The hypothesis of locality is clearly an approximation in the same sense that a curve can be locally approximated by its tangent line. An accelerated observer passes through a continuous infinity of momentarily comoving inertial observers; therefore, the measurements of the accelerated observer could in general be related to the measurements of the whole sequence of hypothetical inertial observers. Let  $\psi$  be the field according to the standard inertial observers and  $\Psi' = \Lambda\psi$  be the field as measured by the instantaneously comoving inertial observer. Specialized to the electromagnetic case,  $\psi$  and  $\Psi'$  are column 6-vectors representing  $F_{\mu\nu}$  and  $F_{(\alpha)(\beta)}$ , respectively, so that  $\Lambda$  is a  $6 \times 6$  matrix determined via equation (1). The most general relationship between the result of field measurement by an accelerated observer  $\Psi$  and  $\Psi'$  that is consistent with causality and the superposition principle is

$$\Psi(\tau) = \Psi'(\tau) + \int_{\tau_0}^{\tau} K(\tau, x)\Psi'(x) dx \quad , \quad (6)$$

where  $\tau_0$  is the instant at which the acceleration is turned on and  $K(x, y)$  is a kernel that must be determined on the basis of further physical hypotheses. If  $K = 0$ , then the hypothesis of locality is satisfied; therefore, we expect the integral part in (6) that signifies the deviation from locality to be of order  $\lambda/\mathcal{L}$ , where  $\lambda$  is the intrinsic scale of fluctuations of the field and  $\mathcal{L}$  is an acceleration length of the observer.

The general theory of integral relationships of the form (6) was originally developed by Vito Volterra [10], who showed that in the space of continuous functions the relationship between  $\Psi$  and  $\Psi'$  is unique. This result has been essentially extended to the Hilbert space of square-integrable functions by Tricomi [11].

Searching for a basic physical hypothesis that would determine the kernel  $K$ , we turn to the phenomenon of spin-rotation coupling. The formula (5),

that is based on the hypothesis of locality for the field, has the simple consequence that the RCP radiation field becomes static ( $\omega' = 0$ ) according to rotating observers for  $\omega = \Omega$ . That is, by a mere rotation an observer can stay completely at rest with a radiation field. Let us recall that if Maxwell's equations are assumed to hold in all inertial frames, then the speed of light  $c$  must be constant for all inertial observers. To ensure this fact, no inertial observer can move with speed  $c$ ; therefore, in the relativistic Doppler formula,  $\omega' = \gamma\omega(1 - \mathbf{v} \cdot \hat{\mathbf{k}}/c)$ ,  $\omega'$  can never be zero. For the motion of an inertial observer along the direction of propagation of the wave,  $\omega' = \omega(1 - \beta)^{1/2}/(1 + \beta)^{1/2}$  can be made as small as possible for  $\beta \rightarrow 1$ , yet the mathematical limit is avoided due to the physical restriction that  $\beta < 1$ . The only case for which  $\omega' = 0$  is that  $\omega = 0$ , i.e. if one inertial observer measures a time-independent field, then all inertial observers measure time-independent fields. In this way the existence of a quantum of radiation as well as the number of quanta in the field acquires an observer-independent status. It is worthwhile to generalize this basic consequence of Lorentz invariance for inertial observers to all observers. Therefore, we demand that if  $\Psi$  becomes constant, then  $\psi$  should be constant as well. Let us note that this basic assumption is related to a natural generalization of quantum mechanics for noninertial observers. A postulate of classical mechanics [12] is that an observer can be comoving with a particle, since they can have the same position and velocity. A classical observer cannot be comoving with a quantum particle, however, as a consequence of Heisenberg's uncertainty principle. The state of the quantum particle in fact satisfies a wave equation and hence a classical observer cannot stay at rest with respect to a fundamental wave.

Before implementing the general requirement that no observer can stay completely at rest with respect to a radiation field, we need to discuss the resolvent kernel for (6). This is done in the next section.

#### 4 Resolvent Kernel

Let us start with an integral equation of the form

$$\phi(x) = f(x) + \lambda_0 \int_a^x K(x, y) \phi(y) dy \quad , \quad (7)$$

where  $\lambda_0$  is a constant parameter. It can be shown that

$$f(x) = \phi(x) + \lambda_0 \int_a^x R(x, y) f(y) dy \quad , \quad (8)$$

where  $R$  is the resolvent kernel [10, 11]. To find  $R$  in terms of  $K$ , let us define the successive iterated kernels of  $K$  by  $K_1(x, z) = K(x, z)$  and

$$K_{n+1}(x, z) = \int_z^x K(x, y)K_n(y, z)dy . \quad (9)$$

It can be shown [10,11] that

$$R(x, y) = - \sum_{n=1}^{\infty} \lambda_0^{n-1} K_n(x, y) . \quad (10)$$

Consider now the special case in which  $K(x, y) = k(x - y)$ , i.e. the kernel is of the convolution (Faltung) type. One can easily show, by letting  $x - y = u$  and  $x - z = t$  in (9), that all the iterated kernels are of convolution type with

$$k_{n+1}(t) = \int_0^t k(u)k_n(t - u)du , \quad (11)$$

i.e. the iterated kernels can be obtained by successive convolutions of  $k$  with itself. Denoting the convolution operation by a star,

$$\phi * \chi(t) = \int_0^t \phi(u)\chi(t - u)du = \chi * \phi(t) , \quad (12)$$

and writing  $\phi * \phi = \phi^{*2}$ , etc., we can express the resolvent kernel  $R(x, y) = r(x - y)$  as

$$r(t) = - \sum_{n=1}^{\infty} \lambda_0^{n-1} k^{*n}(t) . \quad (13)$$

Finally, let us note that if  $K(x, y) = k_0(y)$ , then the iterated kernels with  $n > 1$  and the resolvent kernel are in general functions of both variables.

## 5 Kernel K

The integral equation (6) may be written in the form

$$\Psi(\tau) = \Lambda(\tau)\psi(\tau) + \int_{\tau_0}^{\tau} K(\tau, \tau')\Lambda(\tau')\psi(\tau')d\tau' , \quad (14)$$

such that at  $\tau = \tau_0$ ,  $\Psi(\tau_0) = \Lambda(\tau_0)\psi(\tau_0)$ . To implement our basic hypothesis, we assume that  $\Psi(\tau)$  is constant, i.e.  $\Psi(\tau) = \Psi(\tau_0)$ , and expect that  $\psi(\tau)$  will be constant as well, i.e.  $\psi(\tau) = \psi(\tau_0)$ ; hence,

$$\Lambda(\tau_0) = \Lambda(\tau) + \int_{\tau_0}^{\tau} K(\tau, \tau')\Lambda(\tau')d\tau' . \quad (15)$$

This is the basic integral equation for the determination of  $K$ . Using the notion of the resolvent kernel discussed in the previous section, we can write

$$\Lambda(\tau) = \Lambda(\tau_0) + \int_{\tau_0}^{\tau} R(\tau, \tau') \Lambda(\tau_0) d\tau' , \quad (16)$$

so that we have

$$\int_{\tau_0}^{\tau} R(\tau, \tau') d\tau' = \Lambda(\tau) \Lambda^{-1}(\tau_0) - 1 . \quad (17)$$

This integral relation, in which the right side is known, is not sufficient to determine the resolvent kernel. We need an extra assumption. In this paper we explore two possibilities regarding  $K$ : case (i)  $K(\tau, \tau') = k(\tau - \tau')$  and case (ii)  $K(\tau, \tau') = k_0(\tau')$ .

In the first case, a convolution-type kernel implies that  $R(\tau, \tau') = r(\tau - \tau')$  and (17) becomes

$$\int_0^{\tau - \tau_0} r(u) du = \Lambda(\tau) \Lambda^{-1}(\tau_0) - 1 \quad (18)$$

and a simple differentiation results in

$$r(x) = \frac{d\Lambda(x + \tau_0)}{dx} \Lambda^{-1}(\tau_0) . \quad (19)$$

Once  $r(x)$  is determined from (19), (13) with  $\lambda_0 = -1$  implies, via the reciprocity between  $K$  and  $R$ , that  $K(\tau, \tau') = k(\tau - \tau')$  is given by

$$k(x) = \sum_{n=1}^{\infty} (-1)^n r^{*n}(x) . \quad (20)$$

In particular if  $r(x) = 0$ , then  $k(x) = 0$  and the nonlocality disappears.

In the second case, we have via differentiation of (15) that

$$k_0(\tau) = -\frac{d\Lambda(\tau)}{d\tau} \Lambda^{-1}(\tau) . \quad (21)$$

It is interesting to illustrate these results for the electrodynamics of linearly accelerated systems. Consider an observer at rest on the  $z$ -axis for  $-\infty < \tau < \tau_0$ . At  $\tau = \tau_0$ , the observer accelerates linearly from rest along the  $z$ -axis with acceleration  $g(\tau) > 0$ . Let

$$\theta(\tau) = \int_{\tau_0}^{\tau} g(\tau') d\tau' , \quad (22)$$



$C = \cosh \theta$ , and  $S = \sinh \theta$ ; then, the nonrotating orthonormal tetrad along the observer worldline is given by

$$\begin{aligned}\lambda^\mu_{(0)} &= (C, 0, 0, S) , & \lambda^\mu_{(1)} &= (0, 1, 0, 0) , \\ \lambda^\mu_{(2)} &= (0, 0, 1, 0) , & \lambda^\mu_{(3)} &= (S, 0, 0, C) .\end{aligned}\quad (23)$$

Equation (1) can be written as  $F' = \Lambda F$ , where  $F$  is a column 6-vector with  $\mathbf{E}$  and  $\mathbf{B}$  as components and

$$\Lambda = \begin{bmatrix} U & V \\ -V & U \end{bmatrix} , \quad (24)$$

where

$$U = \begin{bmatrix} C & 0 & 0 \\ 0 & C & 0 \\ 0 & 0 & 1 \end{bmatrix} , \quad V = SI_3 . \quad (25)$$

Here  $I_i, (I_i)_{jk} = -\epsilon_{ijk}$ , is a  $3 \times 3$  matrix that is proportional to the operator of infinitesimal rotations about the  $x^i$ -axis. Let us note that  $\Lambda(\tau_0)$  is the identity matrix in this case.

For the kernel of convolution type, (19) implies that

$$r(\tau - \tau_0) = g(\tau) \begin{bmatrix} R_1 & R_2 \\ -R_2 & R_1 \end{bmatrix} , \quad (26)$$

where  $R_1 = SJ_3, R_2 = CI_3$ , and  $(J_k)_{ij} = \delta_{ij} - \delta_{ik}\delta_{jk}$ . The convolution kernel  $k(x)$  constructed from  $r(x)$  as in (20) is rather complicated in general; however, it takes a very simple form for *uniform* acceleration. In fact,  $k$  is in general a constant for arbitrary *uniform* acceleration. In the case under consideration, an explicit calculation using (20) results in

$$k = -g_0 \begin{bmatrix} 0 & I_3 \\ -I_3 & 0 \end{bmatrix} , \quad (27)$$

where  $g(\tau) = g_0$  is the constant magnitude of acceleration.

On the other hand, for the case  $K(\tau, \tau') = k_0(\tau')$ ,  $k_0$  can be simply computed using (21) and the result is

$$k_0(\tau) = -g(\tau) \begin{bmatrix} 0 & I_3 \\ -I_3 & 0 \end{bmatrix} . \quad (28)$$

Let us remark that for *uniform* acceleration, we have the same result as in the convolution case (27); in fact, this is generally the case for arbitrary constant

acceleration. That is, it turns out that the kernel  $K$  is constant for uniform translational and rotational accelerations and that this unique result has an interesting interpretation in terms of the spacetime connection [13]. However, for nonuniform acceleration the two approaches give different results.

An important distinction between the two approaches is the nature of the memory of nonuniform acceleration that lingers after the acceleration has been turned off, say, at  $\tau_1 > \tau_0$ . In case (i), the field as measured by the observer would still be nonlocal in general, while in case (ii) the field would be local and the integral term in (14) would simply be a *constant* field that can be easily eliminated by a constant recalibration of the measuring devices carried by the observer. Future experiments may use such differences to distinguish between the two possibilities; that is, one could determine whether the kernel  $K$  is of convolution type as in case (i) or  $K(\tau, \tau') = k_0(\tau')$  as in case (ii). In this connection, it is interesting to note here that convolution-type kernels have been employed in the phenomenological studies of material media for a long time. Such studies of history-dependent phenomena apparently began with Poisson's work [14] and have continued in hysteresis theory and continuum physics to the present time [15]. For this reason, a convolution-type kernel was assumed at the outset in the nonlocal theory of accelerated systems presented in [16].

Let us now discuss certain consequences of nonlocality that follow from the theory developed here. It follows from (19) and (21) that for constant  $\Lambda$  the nonlocality disappears. This is the case for all fields if the observer is inertial. Moreover, for a fundamental scalar (or pseudoscalar) field  $\Lambda = 1$  and hence nonlocality disappears. Thus it would be possible in principle for an observer to stay completely at rest with respect to a scalar (or pseudoscalar) field. This is forbidden by our postulate, however. The theory developed here therefore excludes the possibility of existence of a basic scalar (or pseudoscalar) field in nature. Only composite fields of this type can occur; hence, should the Higgs boson be discovered, it would have to be a composite particle.

In connection with the observational consequences of nonlocality, it is interesting to return to the phenomenon of spin-rotation coupling discussed in section 2. Dealing with the uniformly rotating observer, the two cases considered in (19) and (21) result in the same kernel that has been studied in some detail in [16]. A direct consequence of nonlocality in this case is that the field amplitude as measured by the rotating observer depends on the helicity of the incident radiation. According to the hypothesis of locality, the measured amplitudes will be independent of the helicity of perpendicularly incident plane wave, i.e. if the amplitudes of RCP and LCP waves are equal according to the standard observers in the inertial frame, they will also be equal according

to the comoving inertial observers. This has a direct analog in the impulse approximation of quantum scattering theory [17]. On the other hand, nonlocality implies that for  $\omega > \Omega$  the amplitude will be larger by  $(1 - \Omega/\omega)^{-1}$  for RCP radiation and smaller by  $(1 + \Omega/\omega)^{-1}$  for LCP radiation. For microwaves of  $\lambda \simeq 1$  cm incident normally at an observer rotating with a frequency of 300 Hz,  $\Omega/\omega \simeq 10^{-8}$ .

## 6 Discussion

It is important to subject the nonlocal theory developed here to direct experimental test. For this purpose one could re-examine the whole observational basis of the theory of relativity — since such data generally involve observers that are accelerated — to search for evidence of nonlocality. In this connection, an important difficulty is that the electrodynamics of accelerated media is rather poorly understood. It would therefore be a complicated matter to isolate the vacuum nonlocality under consideration here.

On the other hand, if the idea of nonlocality has merit, then we expect that the nonlocal theory would be in better agreement than the standard theory with quantum mechanics in the correspondence limit. For instance, one may consider the accelerated observer to be — in a certain approximate sense — an electron in a Rydberg state. By studying various transition rates, one may be able to provide evidence for the nonlocal theory. In particular, for the uniformly rotating observer, nonlocality implies that there is a relative increase (decrease) in the amplitude of the measured field by  $\Omega/\omega$  ( $-\Omega/\omega$ ) for incident RCP(LCP) radiation with  $\omega \gg \Omega$ . Preliminary efforts in this direction indicate that this prediction of the nonlocal theory is in qualitative agreement with quantum-mechanical results [18].

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