A NEW, NONLOCAL QUANTUM ELECTRODYNAMICS AND ITS EFFECTS ON ATOMIC PHYSICS, COSMOLOGY AND PARTICLE PHYSICS: TIME'S ARROW AND NONLOCALITY IN QUANTUM MECHANICS

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Abstract

We show in this article that as a consequence of a recently demonstrated symmetry of matter fields under a new kind of *nonlocal gauge transformation*, a nonlocal quantum electrodynamics (QED) can be formulated. It is shown that, in the lowest order, this new QED allows for coherent emission/absorption of *two phase-correlated photons* by an atomic system under suitable conditions which can be met in the laboratory as well as in cosmological environments. An experimental test for this two-photon absorption is suggested.

The impact of this result can be felt in many diverse areas in physics, including atomic and molecular physics, cosmology and particle physics. The work also addresses basic questions such as quantum-mechanical nonlocality and a universal time's arrow at the quantum level, and comes up with answers which hinge on a single basic principle — energy conservation at the atomic level.

In particular, interaction of matter with this nonlocal radiation field and with squeezed vacuum are shown to be describable by almost identical mathematics. Several aspects of matter-(squeezed vacuum) interaction as well as multiphoton ionization of matter placed in strong laser fields are discussed and shown to follow as natural consequences of matter-radiation interaction within this nonlocal QED. Furthermore, coherent two-photon absorption from the cosmic microwave background radiation (CMBR) by interstellar matter has been shown to be a viable process, and an effect has been predicted which leads to a clue to the *anisotropy* in the CMBR.

Our conventional wisdom is silent on the question of unidirectionality of the flow of time at the quantum level. We show that in interaction of matter with this nonlocal radiation, a necessary condition for energy conservation at the atomic level is that the present should be determined by the past only and not by the future, which establishes a time's arrow. It is argued that in view of the abovementioned interaction between interstellar matter and the CMBR, the latter can provide us with a *universal* arrow of time. We also show that in situations involving such light-matter interactions, an Einstein-Podolsky-Rosen (EPR)-type nonlocality arises naturally in quantum mechanics as a necessary condition for energy conservation. The nature of the nonlocal correlation is discussed, and it is shown that the nonlocal "link" between two space-time events (which can be spatially separated) is due to a *phase coherence* and has nothing to do with any cause-and-effect chain, thus satisfying the special relativity requirement that two spatially separated events cannot be causally linked.

We also show that Feynman's picture of antiparticles as moving *backward in time* follows as a natural consequence of the fact that in reactions such as particle-antiparticle collisions, pair production, particle decay etc., mass is converted into energy and vice versa. Our work also indicates that the CPT theorem can be extended to nonlocal gauge fields.

I. Introduction and preview

Local gauge invariance has turned out to be a powerful guiding principle in the development of modern physics, the Standard Model (SM) being its most successful achievement. As is well-known, the SM can be symbolically represented as the product of three gauge groups : U(1) $\otimes$ SU(2) $\otimes$ SU(3).

In this review we present some recent results that have been obtained in quantum electrodynamics, the first group in the product above, by replacing the local gauge transformation (GT) in

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$U(1)$ with a *non-local* gauge transformation (Rai Dastidar and Rai Dastidar 1998). We show that the Lagrangian of matter fields remains covariant under this nonlocal GT if, similarly as in the case of elementary quantum field theory, a nonlocal dynamical field is introduced to construct the covariant derivative. This gives a new, nonlocal QED which, *in the first order*, leads to creation and annihilation operators that create/absorb *two correlated photons*. This leads to a host of new results in atomic and molecular physics as well as in astrophysics.

When we impose the requirement of energy conservation in matter-radiation interaction in this nonlocal QED, we obtain two further new results : first, the principle of causality emerges as a necessary condition for energy conservation, thus defining an arrow of time — and secondly, a feature of EPR-type nonlocality in the matter-radiation interaction emerges, also as a necessary condition for energy conservation.

The plan of this review is as follows. In section II we present the basic theory for the nonlocal GT, yielding a nonlocal electromagnetic field which, when inserted into the Lagrangian, meets the requirement of gauge covariance. In section III we work out the interaction of this nonlocal electromagnetic field with matter and obtain all the results referred to above. Finally, in section IV, we summarize our conclusions.

II. Theory

We first establish our notation. We define a nonlocal operator $P|x\rangle\langle x'|$ operating on a local field $\phi(x)$ by the relation $[x = (\vec{r}, t), x' = (\vec{r}', t')]$

$$P\phi = P|x\rangle\langle x'|\phi \equiv \int P(x, x')\phi(x')\, d^4x'$$

(Note that $|x\rangle$ and $\langle x'|$ are not Dirac bra and ket vectors.) With another nonlocal operator $Q|x\rangle\langle x'|$,

$$QP\phi = \int\int Q(x, x')P(x', x'')\phi(x'')\, d^4x'\, d^4x''$$

We remember that in U(1), an arbitrary field $\phi(x)$ transforms[1] like

$$\phi \to \exp\{-ie\chi(x)\}\phi \tag{1}$$

under a local gauge transformation. We propose a *non-local* gauge transformation of the form

$$\phi \to \exp\{-ie\Lambda|x\rangle\langle x'|\}\phi = \int \exp\{-ie\Lambda(x, x')\}\phi(x')\, d^4x' \tag{1'}$$

This can immediately confront the reader with a question : are x and x' within each other's "light cone" ? In other words, does special relativity restrain the range of integration in (1′) ? We shall see later that we are guided to an unambiguous answer to this question very simply and directly by the principle of energy conservation.

We will show now that, as in the case of local gauge transformation in quantum field theory, the coherence property of the lagrangian density of a complex scalar field, is maintained under this new non-local gauge transformation, *provided we replace the ordinary electromagnetic four-potential $A_\mu(x)$ by a non-local potential $\mathcal{A}_\mu|x\rangle\langle x'|$*, which transforms under the above non-local GT as

$$\mathcal{A}_\mu|x\rangle\langle x'| \to \mathcal{A}_\mu|x\rangle\langle x'| + \partial_\mu\Lambda|x\rangle\langle x'| + d_\mu\Lambda|x\rangle\langle x'|, \tag{2}$$

where $\partial_\mu \equiv \partial/\partial x^\mu$, $d_\mu \equiv \partial/\partial x'^\mu$.

[1] We are using units such that $\hbar = c = 1$.

The change in ϕ under this non-local gauge transformation is given by (for infinitesimal $\Lambda(x,x')$) $\delta\phi = -ie\int\Lambda(x,x')\phi(x')\,d^4x'$. One at once notes the parallelism with the Lippmann-Schwinger equation in the first order :

$$\begin{aligned}\Psi &= (1+G_0V)\psi \\ &= \psi(\vec{r}) + \int G_0(\vec{r},\vec{r}')V(\vec{r}')\psi(\vec{r}')\,d^3r',\end{aligned}$$

where ψ is an unperturbed state, Ψ is a scattering state, $G_0(\vec{r},\vec{r}')$ is the non-local Green's function $(E-H_0)^{-1}$ and V is the scattering potential.

Similarly, the newly defined derivative

$$\begin{aligned}D_\mu\phi &= \partial_\mu\phi(x) + d_\mu\phi(x) + ie\mathcal{A}_\mu|x\rangle\langle x'|\phi \\ &= \partial_\mu\phi(x) + d_\mu\phi(x) + ie\int\mathcal{A}_\mu(x,x')\phi(x')\,d^4x' \qquad (3)\end{aligned}$$

is covariant, provided $\mathcal{A}$ transforms under this gauge transformation as

$$\mathcal{A}_\mu\phi \to \mathcal{A}_\mu\phi + \int[\partial_\mu\Lambda(x,x')]\phi(x')\,d^4x' + \int[d_\mu\Lambda(x,x')\phi(x')]\,d^4x' \qquad (3')$$

as may be easily checked using eqn. (2). (Of course, for local ϕ, the second term on the right side of $D_\mu\phi$ vanishes.) All the results derived in the earlier papers and the present paper are obtained without assuming any specific functional form of the quantities $\Lambda(x,x')$ and $\mathcal{A}(x,x')$. It can be readily verified that this transformation property is consistent with the requirement of *gauge invariance* of the non-local (electromagnetic) field tensor

$$\mathcal{F}^{\mu\nu}|x\rangle\langle x'| = (\partial^\mu + d^\mu)\mathcal{A}^\nu|x\rangle\langle x'| - (\partial^\nu + d^\nu)\mathcal{A}^\mu|x\rangle\langle x'| \qquad (4)$$

under this new gauge transformation, *provided that*

$$[\partial^\mu, d^\nu] = 0. \qquad (4')$$

One can write down a non-local Lagrangian in the similar manner as in the case of local Lagrangian, only thing is that one will have to replace A_μ by $g_{\mu\nu}\mathcal{A}^\nu|x\rangle\langle x'|$. Similarly the nonlocal interaction Lagrangian can be shown to be invariant under this gauge transformation [1]. We see, therefore, that our non-local potential *is consistent with the basic requirement of gauge covariance of the Lagrangian.*

In the usual simple harmonic oscillator mode expansion of the electromagnetic field potential **A**, an integral number of photons (0,1,2,...) is associated with each mode. All these modes are orthogonal and independent of each other, and since photons obey the Bose statistics, each mode can be occupied by any arbitrary number of photons. However, in situations such as described above where a *phase-correlated behaviour* of photons becomes manifest, one might think whether a different, *correlated* type of mode expansion could better describe the situation, much as a correlated wave-function gives a better representation of, say, the Helium atom than the orbital product wavefunction does. Since there is no force acting between photons, the observed "correlation" between them is not due to any "configuration interaction" as in the case of a multi-electron system, but due either to the presence of certain non-classical features in the radiation field introduced by a non-linear medium (e.g. squeezing), or to correlated interaction of the radiation field with the atomic electrons, or both.

This type of situation arises in two entirely different scenario; (i) in phase-correlated multiphoton excitation of atoms/molecules and (ii) in interaction of phase-squeezed vacuum field with a three-level atom. Consider an intense classical (thermal or laser) light incident on an atom, and let a pair

of photon absorption events occur at two different space-time points $(\mathbf{r}, t)$ and $(\mathbf{r}', t')$, the interval between which is very close to zero, i.e. $|\mathbf{r} - \mathbf{r}'| \sim c|t - t'|$. This situation can arise when, say, the outer electrons of a Rydberg atom absorb a pair of photons within a time $\delta t \ll \frac{1}{\omega}$. For laser intensities of order 10^{10} W/cm^2 or higher, this is commonplace, if we remember that in the visible range, a flux of one photon $a_0^{-2} t_0^{-1}$ corresponds to $\sim 5 \times 10^{14}$ W/cm^2 of intensity (a_0 and t_0 being the atomic units of length and of time respectively). We have here a situation where two events, which may have occurred far apart in space, are correlated in phase; indeed, if the events are spatially separated, the correlation assumes the character of an EPR-like "quantum entanglement".

Consider now a different scenario : radiation from an optical parametric oscillator generating a phase-squeezed vacuum field is absorbed by a three-level atom having a principal quantum number ≥ 5. If the squeezed vacuum is produced by parametric frequency down-conversion, where an incident pump photon has split up into two lower-frequency photons (signal and idler) which constitute a highly correlated photon pair, a two-photon absorption similar as observed in [2] occurs; i.e. two photons in exact (or nearly exact) phase correlation can be absorbed in two events widely separated in space even at low intensities. Although nothing can be said about the specific mathematical forms of these non-local potentials/field strengths, we will find that we can go a long way to predict specific consequences of this general non-local field. In order to quantise this non-local field, we proceed by carrying out a Fourier expansion [3] of this non-local potential over the usual photon modes:

$$\mathcal{A}(x, x') = \sum_{\mathbf{k}_1 \lambda_1} \sum_{\mathbf{k}_2 \lambda_2} \left[C_{12} \hat{\epsilon}_{12}(\hat{r}, \hat{r}') \exp(i\mathbf{k}_1.\mathbf{r} + i\mathbf{k}_2.\mathbf{r}' - i\omega_{\mathbf{k}_1} t - i\omega_{\mathbf{k}_2} t') + \text{c.c.} \right], \tag{6}$$

where we have written $C_{12} \equiv C_{\mathbf{k}_1 \lambda_1 \mathbf{k}_2 \lambda_2}$, and $\hat{\epsilon}_{12} \equiv \hat{\epsilon}_{\mathbf{k}_1 \lambda_1 \mathbf{k}_2 \lambda_2}$. It is to be mentioned here that we cannot fix the direction of the polarisation vector $\hat{\epsilon}(\hat{r}, \hat{r}')$ "by hand" (as in standard QED) without loss of generality; this is at once obvious if we make the dipole approximation in the interaction hamiltonian. In the next section, to derive the interaction hamiltonian for the matter-field interaction, we shall use for $\hat{\epsilon}(\hat{r}, \hat{r}')$ a power series in $\hat{r}$ and $\hat{r}'$ with undetermined coefficients, which is the most general form possible. However, to quantise the field using the Coulomb gauge we impose for the present the condition that

$$(\hat{\mathbf{k}}_1 + \hat{\mathbf{k}}_2).\hat{\epsilon} = 0. \tag{7}$$

This restriction will be removed later.

We note that the non-local electric field in the Coulomb gauge is given by

$$\mathcal{E}(x, x') = -(1/c) \left(\frac{\partial \mathcal{A}}{\partial t} + \frac{\partial \mathcal{A}}{\partial t'} \right) = \frac{i}{c} \sum_{\mathbf{k}_1 \lambda_1} \sum_{\mathbf{k}_2 \lambda_2} (\omega_{\mathbf{k}_1} + \omega_{\mathbf{k}_2}) \left[C_{12} \hat{\epsilon}_{12}(\hat{r}, \hat{r}') f(x, x') - \text{c.c.} \right], \tag{8}$$

where $f(x, x')$ is the phase factor in (6). The expression for the non-local magnetic field can be found in an analogous manner.

To obtain the field energy at a point $(\mathbf{r}, t)$ we need E^2 and B^2 which we define as

$$E^2(\mathbf{r}, t) = \int \int \mathcal{E}(x, x').\mathcal{E}(x, x'')\, dx'\, dx'',$$

where the limits of integration over dt' and dt'' are determined by the choice of interaction (retarded or advanced). However, in the present work r's are of atomic dimensions, and we simply put $t', t'' \leq t$ for the retarded interaction and $t', t'' \geq t$ for the advanced interaction.

An examination of eqns. (6) and (8) shows that in E^2 and B^2, the crossed terms between two modes $\mathbf{k}_1 \lambda_1$ and $\mathbf{k}_2 \lambda_2$ vanish unless the two modes are strongly phase-correlated. Denoting the population fractions of such correlated-pair modes and of the independent modes by a_{II} and a_{I} respectively, use of standard methods in Quantum Electrodynamics leads to the following expression for the field energy in a volume Ω:

$$W = a_{\text{I}}^2 W_{\text{I}} + a_{\text{II}}^2 W_{\text{II}}, \tag{9}$$

where $W_{\rm I}$ is the usual energy summed over single modes, and

$$W_{\rm II} = \tfrac{1}{2}\sum_{\mathbf{k}\lambda}(Y^2_{\mathbf{k}\lambda} + 4\omega^2_{\mathbf{k}} Z^2_{\mathbf{k}\lambda}), \qquad (9'),$$

where $Y_{\mathbf{k}\lambda} = -i\frac{\sqrt{\Omega/\pi}}{2c}2\omega_{\mathbf{k}}(C_{\mathbf{k}\lambda} - C^*_{\mathbf{k}\lambda})$, $Z_{\mathbf{k}\lambda} = \frac{\sqrt{\Omega/\pi}}{2c}(C_{\mathbf{k}\lambda} + C^*_{\mathbf{k}\lambda})$. Each mode subscript here actually stands for a correlated mode pair $(\mathbf{k}_1\lambda_1, \mathbf{k}_2\lambda_2)$, and $2\omega_{\mathbf{k}}$ stands for $\omega_{\mathbf{k}_1} + \omega_{\mathbf{k}_2}$. Y and Z satisfy the canonical equations of motion with $W_{\rm II}$ as the Hamiltonian as in standard QED, and we now quantise this non-local field by requiring that Y and Z be q-numbers obeying the commutation relation

$$[Z_{\mathbf{k}\lambda}, Y_{\mathbf{k}'\lambda'}] = i\hbar\delta_{\mathbf{k}\mathbf{k}'}\delta_{\lambda\lambda'}. \qquad (10)$$

We thus obtain the non-local field energy in terms of a *new pair of creation and annihilation operators* :

$$W_{\rm II} = \tfrac{1}{2}\sum_{\mathbf{k}\lambda}\hbar\omega_{\mathbf{k}}\left(b_{\mathbf{k}\lambda}b^\dagger_{\mathbf{k}\lambda} + b^\dagger_{\mathbf{k}\lambda}b_{\mathbf{k}\lambda}\right), \qquad (9''),$$

$$\text{where}\quad b_{\mathbf{k}\lambda}, b^\dagger_{\mathbf{k}\lambda} = \frac{1}{\sqrt{2\hbar\omega_{\mathbf{k}}}}\left(2\omega_{\mathbf{k}}Z_{\mathbf{k}\lambda} \pm iY_{\mathbf{k}\lambda}\right). \qquad (11)$$

Equation $(9'')$ is identical in form with standard QED, but the creation and annihilation operators obey a *new commutation relation* :

$$\left[b_{\mathbf{k}\lambda}, b^\dagger_{\mathbf{k}'\lambda'}\right] = 2\delta_{\mathbf{k}\mathbf{k}'}\delta_{\lambda\lambda'}, \quad [b_{\mathbf{k}\lambda}, b_{\mathbf{k}'\lambda'}] = \left[b^\dagger_{\mathbf{k}\lambda}, b^\dagger_{\mathbf{k}'\lambda'}\right] = 0. \qquad (12)$$

as may be directly verified from (10) and (11). The following, and only the following, results of operation of these new correlated-pair mode operators are consistent with the commutation rules (12) :

$$b_{\mathbf{k}\lambda}|n_{\mathbf{k}\lambda}\rangle = \sqrt{n_{\mathbf{k}\lambda}}|n_{\mathbf{k}\lambda} - 2\rangle, \quad b^\dagger_{\mathbf{k}\lambda}|n_{\mathbf{k}\lambda}\rangle = \sqrt{n_{\mathbf{k}\lambda}+2}|n_{\mathbf{k}\lambda}+2\rangle, \quad b^\dagger_{\mathbf{k}\lambda}b_{\mathbf{k}\lambda}|n_{\mathbf{k}\lambda}\rangle = n_{\mathbf{k}\lambda}|n_{\mathbf{k}\lambda}\rangle$$

This shows that b and $b^\dagger$ are *two-photon annihilation/creation operators*, and $b^\dagger b$ is the usual number operator. Once again, by creation of *two photons* we mean the creation of *one photon in each member of a correlated mode pair*, and similarly for annihilation. We can now write down the potential and field strengths in terms of these new operators; thus

$$\mathcal{A}(x,x') = \sum_{\mathbf{k}\lambda}\sqrt{\frac{2\pi\hbar c^2}{\Omega\omega_{\mathbf{k}}}}\hat{\epsilon}_{\mathbf{k}\lambda}(\hat{r},\hat{r}')\left[b_{\mathbf{k}\lambda}f(x,x') + b^\dagger_{\mathbf{k}\lambda}f^*(x,x')\right], \qquad (13a)$$

$$\mathcal{E}(x,x') = i\sum_{\mathbf{k}\lambda}\sqrt{\frac{2\pi\hbar\omega_{\mathbf{k}}}{\Omega}}\hat{\epsilon}_{\mathbf{k}\lambda}(\hat{r},\hat{r}')\left[b_{\mathbf{k}\lambda}f(x,x') - b^\dagger_{\mathbf{k}\lambda}f^*(x,x')\right]. \qquad (13b)$$

III. Interaction of non-local radiation field with matter

We show below that the non-local model of QED as expounded above serves to bring out the feature of (phase) correlation between two (or more) photon modes. Consider two zero-order dressed states $|\Psi_f\rangle = |\psi_f\rangle|n_f\rangle$ and $|\Psi_i\rangle = |\psi_i\rangle|n_i\rangle$. In the Coulomb gauge the radiation interaction matrix element coupling these two states is $M_{fi}(t) = -\frac{e}{mc}\langle\psi_f(\mathbf{r}), n_f|\mathbf{p}.\mathbf{A}(\mathbf{r},t)|\psi_i(\mathbf{r}), n_i\rangle$. Here we introduce a non-local potential $\mathcal{A}(\mathbf{r},t,\mathbf{r}',t')$ such that the above matrix element becomes

$$M_{fi}(t) = -\frac{e}{mc}\langle\psi_f(x), n_f|\mathbf{p}.\mathcal{A}(x,x')|\psi_i(x'), n_i,\rangle, \qquad \text{where}\quad x \equiv (\mathbf{r},t), \quad x' \equiv (\mathbf{r}',t'), \qquad (14)$$

i.e. the interaction at the point x is correlated with the interactions at all other points x'. This non-local interaction had been defined in [4] in an *ad hoc* manner; in this review we have shown that, how such a potential can be introduced field-theoretically by demanding covariance under a non-local U(1) gauge transformation.

The integration in (14) runs over $d\mathbf{r}$ and dx', and in this paper we set the limits of the time integral dt' either (i) from $-\infty$ to $t-\frac{|\mathbf{r}-\mathbf{r}'|}{c}$ (giving the retarded interaction), or (ii) from $t+\frac{|\mathbf{r}-\mathbf{r}'|}{c}$ to ∞ (which gives the advanced interaction), leaving out EPR "quantum entanglement"-like correlations for the present. We shall discuss the question of causality later and shall find that this restriction turns out to be unnecessary.

We can equivalently use a similar non-local form of the multipolar interaction at the space-time point x:

$$-e\int \mathbf{r}.\mathcal{E}(x,x')\,dx', \tag{14'}$$

The results obtained so far from our non-locality postulate may be summed up as follows: quantization of a radiation field with a high degree of second-order coherence can depart from the usual QED and can lead to photon modes (more exactly: correlated mode pairs) generated by two-photon creation and annihilation operators. Indeed, if we let a zero-order dressed state of an atom $|\psi\rangle|n\rangle$ evolve in the presence of the non- local electron-radiation interaction (14) or (14′), we obtain a linear superposition of dressed states of the form

$$|\psi\rangle\sum_{k=0}^{\infty}c_k(t)|2k\rangle \tag{15}$$

for even n ; for odd n, $|2k\rangle$ is replaced by $|2k+1\rangle$. The possibility of emission/absorption of two photons in the first order of interaction, i.e. linear in field intensity becomes evident at once as a consequence of this model. Also, the commutation relations (12) are gauge-independent, and hence from now on we can do away with the transversality restriction as mentioned after eqn.(7). Before we work out the possible consequences of this non-local field in a typical atomic transition, let us note its formal similarity with a *squeezed* state. We recall the definition of the squeezed state [5]:

$$|\alpha,\zeta\rangle = D(\alpha)S(\zeta)|0\rangle, \tag{16}$$

where the *squeeze operator* $S(\zeta)=\exp[\frac{1}{2}\zeta^*a^2-\frac{1}{2}\zeta(a^\dagger)^2]$, and the *displacement operator* $D(\alpha)=\exp\left(\alpha a^\dagger-\alpha^*a\right)=e^{\frac{-|\alpha|^2}{2}}e^{\alpha a^\dagger}e^{-\alpha^*a}$. ($\alpha$ and ζ are arbitrary complex numbers.) Putting $\alpha=0$ gives the squeezed vacuum, while putting $\zeta=0$ gives the coherent state. An expression for the squeezed vacuum in terms of the number states has been given by Hollenhorst [6]; putting $\zeta=re^{i\theta}$, this becomes

$$S(\zeta)|0\rangle=(\cosh r)^{\frac{1}{2}}\sum_{n=0}^{\infty}\left(\tfrac{1}{2}e^{i\theta}\tanh r\right)^n\frac{\sqrt{2n!}}{n!}\,|2n\rangle. \tag{17}$$

Thus we see that the squeezed vacuum $S(\zeta)|0\rangle$ has the same form as the expansion (15) above. The *two-mode squeezed vacuum* [7], where the squeezing operator is given by the more generalised form

$$S(\zeta)=\exp[\zeta(\omega)a^\dagger(\omega)a^\dagger(2\Omega-\omega)-\zeta^*(\omega)a(\omega)a(2\Omega-\omega)]$$

provides another example where the doubly occupied modes in (17) would be replaced by singly occupied correlated mode pairs; the signal and idler photons generated in a parametric down-conversion present themselves at once as forming such a correlated pair. We can therefore say that the squeezing operator $S(\zeta)$ is mathematically identical (except for coefficients/multiplying factors) to the evolution operator in a non-local field, and hence atoms interacting with a squeezed field and

with an intense radiation field (for which a non-local description is valid) can be expected to show the same non-classical features. We find below that it is indeed so.

We now work out two typical first-order transition matrix elements using this non-local description for a radiation field. First it must be mentioned that depending on the photon flux and on the coherence properties of the laser field, one must in general describe the field as

$$\mathcal{E} = a_L \mathbf{E}_L + a_N \mathcal{E}_N, \tag{18}$$

where the subscripts L and N refer to the local (usual) and the non-local fields respectively (a_L and a_N are identical with $a_{\rm I}$ and $a_{\rm II}$ in eqn.(9)); the relative strength of the two coefficients would depend, essentially, upon the degree to which the response of the atomic electron/s to the field appears phase-correlated.

We work within the dipole approximation, and formulate the typical matrix element of this two-photon absorption operator (I) for a single-electron atom, and (II) for a two-electron atom.

(I) For simplicity we confine ourselves to a single radiation mode $|n_{\mathbf{k}\lambda}\rangle$. Define

$$T_{\rm I} = \langle \psi_{l_f m_f}(x), n_{\mathbf{k}\lambda} - 2 | \mathbf{r}.\mathcal{E}_N(x,x') | \psi_{l_i,m_i}(x'), n_{\mathbf{k}\lambda} \rangle, \tag{19}$$

where the l's are the orbital angular momenta in the initial and final states. We use equation (13b) for $\mathcal{E}_N(x,x')$ and note that the time integration in (19) for the non-local interaction yields the required energy balance $E_f = E_i + 2\hbar\omega_{\mathbf{k}}$ *if and only if we use the retarded interaction* ; use of the advanced interaction fails to satisfy energy conservation. Let us go into the mathematical details here. We will show now that only the retarded interaction can lead to energy conservation. The time-integral in (19), with retarded correlation, is given by

$$\int_{-\infty}^{\infty} e^{\frac{i}{\hbar}(E_f t - \hbar\omega t)}\, dt \int_{-\infty}^{t-\rho/c} e^{-\frac{i}{\hbar}(E_i t' + \hbar\omega t')}\, dt' = \int_{-\infty}^{\infty} f(t)\, dt \int_{-\infty}^{t-\rho/c} g(t')\, dt' \quad \text{(say)} \tag{20}$$

($\rho = |\mathbf{r} - \mathbf{r}'|$), while with the advanced correlation we have

$$\int_{-\infty}^{\infty} f(t)\, dt \int_{t+\rho/c}^{\infty} g(t')\, dt', \tag{21}$$

The expression (20) can be written as a sum of two integrals :

$$\int_{-\infty}^{0} f(t)\, dt \int_{-\infty}^{t-\rho/c} g(t')\, dt' + \int_{0}^{\infty} f(t)\, dt \int_{-\infty}^{t-\rho/c} g(t')\, dt' \tag{20'}$$

while the expression (21) can be written as

$$\int_{-\infty}^{0} f(t)\, dt \int_{t+\rho/c}^{\infty} g(t')\, dt' + \int_{0}^{\infty} f(t)\, dt \int_{t+\rho/c}^{\infty} g(t')\, dt'. \tag{21'}$$

Elementary considerations show that the first integral in (20') and the second integral in (21') contribute nothing, while the other two integrals are (to within factors) given by $\delta(E_f, E_i + 2\hbar\omega)$. Energy conservation in the two-photon transition thus results, as it were, when we sum up our "samplings" of the phase of the matrix element over extended periods of time.

We now come to the question of the significance of the *limits* of the integrals over dt. The matrix element is, so to say, **created** when we have photoexcited the atom ; let us agree to define that this act has been done at the instant $t = 0$. The crucial point to remember is that although the photoexcitation is not an irreversible process (the atom can certainly be de-excited again), *our acts of exciting and subsequently observing the atom are irreversible in time*. Therefore we can sample

the phase (or whatever else we like to) of the excitation matrix element only at times $t \geq 0$. Thus the first integral in (21'), although non-vanishing, is actually non-physical (it is just as meaningful as sampling the behavioural pattern of a baby before the baby is born), and we are left with only one non-vanishing integral in (20')

$$\int_0^\infty e^{\frac{i}{\hbar}(E_f t - \hbar\omega t)}\, dt \int_{-\infty}^{t-\rho/c} e^{-\frac{i}{\hbar}(E_i t' + \hbar\omega t')}\, dt'. \tag{22}$$

Thus energy conservation will be satisfied if and only if temporal correlation is provided by the past, and not by the future; we have derived an arrow of time in quantum physics as a necessary condition for energy conservation. This remarkable result shows that causality as we understand it — *ability of the past, but not of the future, to influence the present* — follows as a necessary condition for energy conservation in our non-local picture of the electromagnetic field; we have obtained a non-thermodynamic arrow of time having a purely quantum nature.

Incidentally, energy conservation is satisfied even if we set the upper limit of the integration over dt' (eqn.(22)) to t *exactly*, i.e. put $c = \infty$, which corresponds to an EPR-like "quantum entanglement". Thus we find, as had been mentioned in the beginning, that such situations need not be left out from our non-local picture, and that causality in a quantum system need not be restricted by the special theory of relativity. However, if we go further in a closer analysis of the integration over dt', we find that we can make an even stronger statement. Consider the case where $|\mathbf{r} - \mathbf{r}'| > ct$, in which case the upper limit in the t'-integral in (22) becomes negative (remember that both $\mathbf{r}$ and $\mathbf{r}'$ extend over all space), and the integral vanishes, leading to non-conservation of energy ! Thus in our non-local picture of the quantum world, attempts to combine causality with the special theory of relativity are not only unnecessary, but can also lead to results which do not agree with experience; one is immediately led to think of Bell's theorem. Note that such an EPR-like correlation between x and x' is also indicated by the commutation relation (5), which gives the necessary condition for *gauge invariance* of the non-local electromagnetic field tensor.

Finally we get (in atomic units)

$$T_{\mathrm{I}} = i\frac{\sqrt{2\pi I}}{c}\langle\psi_f(\mathbf{r})|\mathbf{r}.\hat{\epsilon}(\hat{r},\hat{r}')|\psi_i(\mathbf{r}')\rangle. \tag{23}$$

As mentioned earlier, we use for $\hat{\epsilon}(\hat{r},\hat{r}')$ a power series in $\hat{r}$ and $\hat{r}'$ with undetermined coefficients; the specific form chosen is

$$\hat{\epsilon}(\hat{r},\hat{r}') = (\hat{r} + \hat{r}')\sum_{n=0}^{\infty} a_n P_n(\hat{r}.\hat{r}') \tag{24}$$

(Symmetry demands that the coefficients of $\hat{r}$ and of $\hat{r}'$ be identical.) Putting $\psi_i(\mathbf{r}) = N_i\frac{R_i(r)}{r}Y_{l_i m_i}(\hat{r})$ and similarly for ψ_f gives

$$T_{\mathrm{I}} = iN_iN_f \times \frac{\sqrt{2\pi I}}{c}\mathcal{R}(\mathcal{I}_1 + \mathcal{I}_2), \tag{25}$$

where $\mathcal{R}$ is a radial integral, and the $\mathcal{I}$'s are two angular integrals, one of which gives the selection rule $l_f = l_i$, while the other gives $l_f = l_i$, $l_i \pm 2$. (See Appendix A for details.) Thus from (25) we obtain two-photon transitions of the type $S \to S$, $S \to D\ldots$ linear in intensity. We also note that the squeezed vacuum corresponds to the limit $\alpha \to 0$ in eqn. (16); an increase in α, which results among other things in an increase in intensity, leads to a deviation of the squeezed field from the form (17) to a situation where the summation over mode indexes runs over all integers (instead of *only even or only odd integers* as in (15)), so that in the first order, an increase in intensity in squeezed light-matter interaction causes, progressively, single-photon transitions instead of the two-photon transitions as was the case in the squeezed vacuum. This result exactly matches the

observation of Georgiades et al [2]. Other possible two-photon absorptions linear in intensity in a squeezed vacuum can be trivially predicted. Two possibilities immediately come to mind that should be experimentally feasible : (i) two-photon ionisation of a Rydberg atom, and (ii) two-photon electron detachment from a negative ion, both with a squeezed vacuum. Indeed, a whole field of non-destructive two-photon spectroscopy of atoms/molecules opens up.

At this point we should remark that the formal similarity between a squeezed field and a laser field with strong second-order coherence can become misleading when we consider the intensity variation of the fields interacting with matter. As we shall soon see, the higher the intensity, the more phase-correlated does the classical (thermal or laser) radiation field "appear" to two detectors (e.g. two atomic electrons), or one detector at two different times, rendering this non-local description the more valid. On the other hand, as we just saw, increase in intensity in a squeezed vacuum takes us back, so to say, into the classical domain (the two-photon transition becomes non-linear with increase in intensity).

(II) For a two-electron atom we require a matrix element of the type

$$T_{\rm II} = \langle \psi_{l_f l'_f}(x,x'), n_{\mathbf{k}\lambda} - 2 | (\mathbf{r}+\mathbf{r}').\mathcal{E}_N(x,x') | \psi_{l_i,l'_i}(x,x'), n_{\mathbf{k}\lambda} \rangle. \tag{26}$$

As before, the time integration in (26) yields the energy balance $\delta(E_f, E_i + 2\hbar\omega_{\mathbf{k}})$, leaving

$$T_{\rm II} = i\frac{\sqrt{2\pi I}}{c} \langle \psi_f(\mathbf{r},\mathbf{r}') | (\mathbf{r}+\mathbf{r}').\hat{\epsilon}(\hat{r},\hat{r}') | \psi_i(\mathbf{r},\mathbf{r}') \rangle \tag{27}$$

To evaluate $T_{\rm II}$ we use a correlated wavefunction of the form [8]

$$\psi = R^{-\frac{5}{2}} \sum_{\mu} F_\mu(R) \Phi_\mu(R;\alpha,\hat{r},\hat{r}'), \tag{28}$$

where $R = \sqrt{r^2 + r'^2}$ and $\alpha = \arctan(r'/r)$ are the hyperspherical coordinates.

Substituting eqns. (28) and (24) in (27) we obtain (see Appendix B for details)

$$\begin{aligned} T_{\rm II} \quad &= i\tfrac{\sqrt{2\pi I}}{c} \int F_f^*(R) R F_i(R)\, dR \int A_f A_i P_{n_f}^{(a_f,b_f)}(\cos 2\alpha)(\cos\alpha)^{l_f+l_i} \\ &\quad \times (\sin\alpha)^{l'_f+l'_i}(\cos\alpha + \sin\alpha) P_{n_i}^{(a_i,b_i)}(\cos 2\alpha) \cos^2\alpha \sin^2\alpha\, d\alpha \\ &\int\int \left[\mathcal{Y}^{M_f}_{L_f l_f l'_f}(\hat{r},\hat{r}')\right]^* (1 + P_1(\hat{r}.\hat{r}') \textstyle\sum_{k=0}^{\infty} a_k P_k(\hat{r}.\hat{r}') \mathcal{Y}^{M_i}_{L_i l_i l'_i}(\hat{r},\hat{r}')\, d\hat{r}\, d\hat{r}' \quad . \end{aligned} \tag{29}$$

We obtain the selection rules for l_f and l'_f from the two angular integrals, one involving $\sum_k a_k P_k(\hat{r}.\hat{r}')$ and the other involving $\sum_k a_k P_1(\hat{r}.\hat{r}') P_k(\hat{r}.\hat{r}')$. They are tedious but straightforward ; note that although the coefficients a_k are unknown, the relative strength of the two terms can be analytically obtained. For the first integral, both Δl and $\Delta l'$ are even for even k and odd for odd k, where as for the second integral Δl and $\Delta l'$ turn out to be odd for even k and even for odd k. Thus finally we see that parity is conserved, as in the single-electron case.

For carrying out the integration over $d\alpha$ one can use the result [8]

$$\begin{aligned} &\int_{-1}^{1} (1-x)^\rho (1+x)^\sigma P_n^{(\alpha,\beta)}(x) P_m^{(\gamma,\delta)}(x)\, dx \\ &= \frac{2^{\rho+\sigma+1}(1+\alpha)_n (1+\gamma)_m \Gamma(\sigma+1)\Gamma(\rho+1)}{m!n!\Gamma(\rho+\sigma+2)} \sum_{r=0}^{m} \frac{(-m)_r (1+\gamma+\delta+m)_r (1+\rho)_r}{r!(1+\gamma)_r (\rho+\sigma+2)_r} \\ &\times {}_3F_2(-n, 1+\alpha+\beta+n, 1+\rho+r; 1+\alpha, 2+\rho+\sigma+r; 1). \end{aligned} \tag{30}$$

On symmetry grounds, $\Delta n = n_f - n_i$ is always even ; sample calculations show that the integral over $d\alpha$ is strongly peaked when $n_f = n_i$. We see therefore that this two-photon transition *maintains the radial correlation* in the two-electron atom.

We can now summarise the properties of the matrix element T as follows :

(i) It is a two-photon excitation linear in the field strength ;

(ii) It conserves the parity of the atom ;

(iii) For a two-electron atom, it conserves the degree of radial correlation.

In view of eqn.(18) this double excitation will be running, in general, parallel to usual single-photon transitions in an experiment, and hence in any multiphoton excitation process in a field having a strong second-order coherence the overall transition would have the form $TTDDTD\ldots$ or some such random mixture (D's are single-photon transition moments), tending towards the extreme form $TTTT\ldots$ as the field grows more intense. If we start out with a closed-shell atom in an S-state, such a series of T-transitions would result in the formation of strongly correlated doubly excited states which are known to favour simultaneous double electron escape (Wannier [8]; see also the review articles of Fano [8] and Rau [8]). This is to be contrasted with *sequential* two-electron escape, which does not involve any electron correlation.

Furthermore for this correlated two-photon absoption by the gas in interaction with the black body radiation, the spectrum of the radiation would appear to us to be slighty deviated from the Planck formula, the amount of the deviation being directly proportional to the density of the gaseous matter. Moreover a non-uniform distribution of this matter would appear to us as an apparent spatial anisotropy of this radiation field temperature.

Here we present the theory for the observed reduction in black body radiation density. The energy density of black body radiation at temperature T (deg K) within a frequency band $\nu \to \nu + d\nu$ is given by the Planck formula

$$E(\nu)d\nu = A(\frac{u^3}{e^u - 1})d\nu, \text{ where } A = \frac{8\pi}{h^2}(\frac{kT}{c})^3, u = \frac{h\nu}{kT}. \tag{31}$$

The function $u^3/(e^u - 1)$ is maximum at $u \approx 2.821439$, i.e. the black body temperature T is related to the frequency ν_m at the peak by T$\approx h\nu_m/(2.821439k)$.

It has been mentioned before that when matter-radiation interaction satisfies the condition of a second order coherence, the symmetry of invariance under a non-local gauge transformation becomes prominent and causes two photon absorption(s) linear in the intensity. As a result, when a black body radiation reaches an observer through a gaseous absorber under circumstances that the above condition is fulfilled, a coherent absorption of two photon of frequency $\frac{1}{2}\nu$ would be superposed as an absorption structure as a frequency ν upon the black body spectrum as viewed by the observer. Under the simplest assumption that the absorption coefficient z is the same over the entire spectral range of interest, the loss $F(\nu)$ as a consequence of the correlated absorption of two photons each of the frequency $\frac{1}{2}\nu$ would be given by

$$F(\nu) = zE(\nu/2) = zA(\frac{y^3}{e^y - 1}), \tag{32}$$

where $y = \frac{1}{2}\nu$. As a result the observed black body radiation density reduces to

$$G(\nu) = E(\nu) - F(\nu) = Af(u, z), \tag{33}$$

where

$$f(u, z) = \frac{u^3}{e^u - 1} - \frac{\frac{1}{8}u^3}{e^{1/2u} - 1}z. \tag{34}$$

IV. Discussions and conclusions

We have given a model for non-local electromagnetic field (derived field theoretically) and will discuss briefly the consequences of the interaction of this non-local field with matter. We will show that the impact of this interaction can be felt in various fields of physics, including atomic and molecular and optical physics, cosmology and particle physics. Some important results are:

(i) As a consequence of matter-(non-local radiation) interaction, a correlated two-photon absorption *linear in intensity* occurs. Direct experimental evidence for such two-photon absorption exists [2]. Many of the features observed in multiphoton double ionization and above- threshold ionization in atoms placed in strong laser fields can be shown to follow as natural consequences of this non-local field model.

(ii) This matter-non-local radiation interaction can also give an explanation for a slight deviation from the Planck formula, in the black body radiation viewed through gaseous matter. We also show that a non-uniformity in this gaseous matter distribution leads to an (apparent) spatial anisotropy of the black body radiation [9].

(iii) The *principle of causality* emerges as a necessary condition for energy conservation in such correlated two-photon absorption processes; out of the entire conceivable range of non-local temporal correlations (which spans both the past and the future), only the past ensures energy conservation in matter-radiation interaction events of the above kind, thus defining a non-thermodynamic arrow of time in the quantum (atomic) level. The cosmic microwave background radiation can provide this time's arrow to all particles in the universe that are acted on by the electromagnetic field [10].

(iv) The self-consistency relation (5), which implies that the two space-time points x, x' are completely independent (i.e. need not lie within the light cone of one another), shows that the non-local correlation must be of the Einstein-Podolsky-Rosen (EPR)-type. It was shown in [1] that this EPR-like character of the non-local correlation in the new QED *is another necessary condition for energy conservation.*

(v) Since the universe consists of only the energy and massive matter, we argue that as a consequence of the requirement for energy conservation in establishing a nonthermodynamic arrow of time mentioned in (iii), particle and anti-particle must necessarily move in opposite directions in time. This result also indicates that the CPT theorem can be extended to cover non-local gauge fields [11].

The purpose of the present review is to show that the cosmic microwave background radiation (CMBR) provides a *via media* for imparting each and every particle (which is acted on by an electromagnetic field) in the universe with this arrow of time. We proceed to demonstrate this by recalling from a recent work [5] that the interaction of matter with CMBR fits into the nonlocal picture as laid out in [2], and thus allows for correlated two-photon absorption by all particles, whereupon the time's arrow is established by energy conservation.

We have shown that if a radiation field (either as generated, or as detected by atomic electrons) enjoys a large second-order coherence — in other words, two photons far apart in space can be *coherently* absorbed by the atomic electrons within a phase difference $\Delta\phi = \omega\delta t \ll 1$ (δt being the time interval between the two photon-absorption events) — then the field is describable by the nonlocal QED. In practice, this translates to the condition that the photon density be significantly higher than the absorber (electron) density. (Of course, if the effect is to be measurable, the electron density must at the same time be high enough to make the observation of the two-photon

absorption process experimentally viable.) Some typical matrix elements for such correlated two-photon absorption by atoms have been worked out in section III. The question is : under what circumstances can this coherence condition be satisfied in practice ? Two suitable laboratory settings have already been described above, namely, (i) a laser field so intense that the average photon density is much larger than the electron density, and (ii) a squeezed light field with phase-correlated photon pairs, e.g. *signal* and *idler* photons. We have also shown [9] that a third, *cosmological* setting is provided by this universe through the *via media* of the cosmic microwave background radiation (CMBR). Indeed, we can call the CMBR the "Lowest common denominator" from the viewpoint of this paper, in the sense that it influences *each particle of the universe capable of interacting with an electromagnetic field*. Note that, in the entire universe, the overall photon density is several orders of magnitude higher than the baryon or lepton density. As such, if we remember that the CMBR *is incident from all directions*, it is obvious that so long as the matter density is tenuous enough (so that the photon density remains higher than the matter density), two distinct photon absorption events can always take place at two arbitrary space points so that the above second-order coherence condition is satisfied. For example, it is easy to see that if two millimetre-wave photons are absorbed within a gap of $\sim 0.1 - 1$ ps, then the phase gap between them is $\ll 1$. The difference between the laboratory settings and the cosmological setting is that, whereas in the former the spatial distance between the two distinct (photon absorption) events is of the order of atomic/molecular dimensions, in the latter the said distance can be of macroscopic dimensions.

Before going into detailed discussions on consequences of correlated two-photon absorption by gaseous medium in interaction with black body radiation, let us demonstrate briefly the observed features in multiphoton ionization in atoms placed in strong laser fields, which can be shown to follow as natural consequences of this non-local field model. As mentioned above, interaction of matter with a strong radiation field fits into this non-local picture if two photons are absorbed by an atom within a time interval $\delta t \ll \frac{1}{\omega}$. For practical purposes this requirement can be replaced by a more convenient one, namely that in a time $\frac{1}{\omega}$, the number of photons flowing into and out of the "volume" V of the atomic shell occupied by the outer electron/s be much larger than the number of these outer electrons. Thus for an n-(outer)-electron atom the condition reads

$$\frac{VI}{c\omega^2} \gg n. \tag{35}$$

On the other hand, when the correlated motion of two electrons is involved (see discussion after eqn.(22)), we have to remember that the laser electric field should be significantly smaller than the Coulomb fields within the atom, in order that the electrons can maintain their mutual correlation. Choosing a field of order 10^{-2} a.u. gives us a crude estimate for an upper limit

$$I_{max} \sim 10^{13} \text{ W/cm}^2 \tag{36}$$

to the laser intensity upto which we can expect electron correlation in a two-electron atom to be maintained.

We now examine some physical consequences of our model QED and compare them with observed features in Multiphoton Ionisation (MPI) and Above-threshold Ionisation (ATI) experiments. Earlier, the occurrence of a TT..-like chain was shown to lead to the possibility of *direct double ionisation* (as opposed to *sequential* double ionisation) [13]; we now see that the possibility can materialise if and only if both the conditions (35) and (36) are fulfilled. Simple considerations show that for short wavelengths the two conditions can become mutually exclusive, thus ruling out the possibility of Direct Double Ionization (DDI), whereas for long wavelengths both conditions (35) and (36) can be satisfied over a broad range of intensities. There is a caveat, though; for *very large* wavelengths the condition (35) can be satisfied at quite low intensities, but the rather large number

of D's in the multiphoton excitation chain $TTDDTD\ldots$ will tend to spoil the electron correlation pattern. Also, for a given power output, the peak intensity in short laser pulses is higher than in longer ones, and hence we should expect DDI to occur more readily with shorter pulses than with long pulses. This accords with the observations of L'Huillier et al [13] who measured multiphoton ionisation in Xenon using both nanosecond and picosecond pulsed lasers at 1064 nm and 532 nm, and those of Agostini and Petite [14] and Delone et al [15] who measured MPI of Ca at 1064 nm with picosecond pulses and nanosecond pulses respectively. Furthermore, Yergeau et al [16] found that in MPI of rare gases using a CO_2 laser ($\lambda = 9.55\mu$ and 10.55μ), the multiply charged ions are formed *sequentially*.

In a strong laser field, another interesting effect stems from the photon number-phase uncertainty relation $\Delta N \Delta\phi \geq 1$. The atomic electron/s during the two-photon excitation is/are acting as a device that measures the incident radiation phase to within an accuracy $\Delta\phi = \omega\delta t$. Consequently the photon number would be uncertain by $\Delta N \sim \frac{1}{\omega\delta t}$, and this would shift the photoelectron spectrum by an amount $\Delta N\hbar\omega$, resulting in the suppression of ΔN lower-order peaks. This is a very well-known phenomenon in ATI (e.g. [17]); attempts to explain it have so far met with only partial success. Note that in the case of two electrons, they can act individually as such phase-measuring devices and no correlation is required; as such, this feature is not restricted by the condition (36). Indeed, peak suppression (also known as quenching) has been observed at very high intensities ($\sim 10^{16}$ W/cm^2).

Use of circularly polarised light lends a particularly interesting twist to this aspect. A circularly polarised photon can be regarded as a superposition of two linearly polarised components 90^o apart in phase, and hence even if two circularly polarised photons arrive with a phase gap $\omega\delta t = \frac{\pi}{2}$ between them, the x-component of one photon and the y-component of the other can still be in exact phase correlation. This implies that the coherence condition $\delta t \ll \frac{1}{\omega}$ can be relaxed to $\delta t \sim \frac{1}{\omega}$, i.e. it would be satisfied at much smaller intensities. Thus at any given intensity, many more lower-order ATI peaks can be expected to be quenched by circularly polarised light than by linearly polarised light. This feature has indeed been observed (e.g. Bucksbaum et al [18], Bashkansky et al [19]). Interestingly, this also allows us to *predict* that with a circularly polarised laser, onset of DDI should occur in MPI experiments at lower intensities than with linearly polarised lasers.

The consequences of such a correlated two-photon absorption loss upon the continuous spectrum of a blackbody radiation travelling through gaseous matter has been worked out in section III. It has been shown that the blackbody radiation energy density $E(\nu)$ at a frequency ν which, in the absence of any absorbing matter, is given by the Planck formula, can be reduced to

$$G(\nu) = E(\nu) - F(\nu), \quad F(\nu) = \alpha E(\nu/2),$$

where α is an effective correlated two-photon absorption factor for the intervening medium. For the present we need not go into the detailed dynamics needed for calculating α; we only note that it will depend, among other things, on the *density of the absorber*. Figure 1 shows the functions $E(\nu)$, $F(\nu)$ and $G(\nu)$ (times 1/A) at a temperature $T = 2.728K$ plotted against ν, with an arbitrarily chosen absorption coefficient $z = 5\%$. It is evident that the peak of the $G(\nu)$-curve, which represents the intensity spectrum viewed by the observer through the gaseous matter, is shifted slightly towards the left of the peak of the PLanck curve $E(\nu)$, and this shift would appear to the observer as an apparent shift ΔT in the radiation temperature - in this particular example the temperature shift turns out to be $-0.039K$, i.e. about 1.4%. Sample calculations with different values of z show that the temperature shift $\Delta T/T$ is of the order of $0.3z$ always.

We have shown above that the theory predicts an apparent shift in the "blackbody temperature" of the radiation travelling through the gas. An experimental test for checking this prediction is being suggested here. For this experiment, one needs a hollow cavity "black body" maintained at any arbitrary temperature, with a spectrophotometer mounted inside and filled with any desired gas,

and record the spectrum at different pressures of the gas. So long as the pressure is sufficiently low so that the electron density is significantly lower than the photon density, the above-mentioned second order coherence condition applies, and the measured black body temperature should show a variation with the gas pressure as given in Eq. (33). Furthermore, this pattern of linear variation of ΔT with z can be expected to hold only so long as the atomic/molecular density of the gaseous matter remains much less than the photon density of the radiation field.

Is there any observational evidence that such a two-photon absorption from the CMBR really does take place anywhere in the universe? There is, such a two-photon absorption loss superimposed upon the spectrum of the CMBR travelling through interstellar and intergalactic matter manifests itself as an apparent deviation of the spectrum from the Planck formula — the magnitude of the deviation depends, of course, on the value of the absorption factor α — and hence this apparent deviation should be more "visible" when the CMBR is observed from within some dense gaseous atmosphere than from within more tenuous atmospheres. We now draw the reader's attention to the balloon-borne measurements of the cosmic microwave spectrum carried out by Woody and Richards [20] and the measurements carried out in the *COBE* (Cosmic Background Explorer) satellite [21], which fit into these two categories admirably. The spectrum observed by Woody and Richards [20] (see Fig. 2 therein) shows a distinct dip below the Planck curve over extended regions towards the right of the spectral peak, the maximum deviation occuring around *twice the peak frequency*, whereas in the *COBE* measurements the deviations from the Planck curve are much more miniscule.

To avoid misgivings that we are jumping to conclusions, we emphasize two points. First, whether such a deviation from the Planck distribution as predicted by our theory does occur at all is a question that can only be settled by experiment; we have suggested one possible experimental test: Secondly, the *viability* of such a correlated two-photon absorption process under conditions of Woody and Richards' balloon measurements has also been established in [9]. Therefore, subject to the condition that our prediction is experimentally verified, we can claim that the CMBR travelling through tenuous matter can satisfy the condition for second-order coherence and hence, as laid down by the non-local QED, is capable of giving rise to *correlated two-photon absorption throughout the universe* — thus offering a very simple picture of the arrow of time in the entire universe at the quantum (atomic) level *as a necessary condition for energy conservation*.

Our picture does not depend on any particular cosmological model. However, if one couples it to the Big Bang model, it would appear that the time's arrow (and hence *causality*) were set in motion as soon as the universe started off — the latter "creating", as it were, its own flow of time as well as its own physical laws as soon as it was born. Appealing as the picture is, it is still incomplete; until the non-local gauge symmetry of the universe can be extended to cover the entire range of electroweak interactions, our theory remains confined only to those constituents of the universe that are acted on by the electromagnetic field and hence cannot embrace neutrinos, which are generally supposed to make up the *cold dark matter* in the universe. So far we have considered phenomena which satisfy energy conservation. Let us now consider reactions which do not conserve energy, but involve mass$\leftrightarrow$energy conversion, e.g. particle-antiparticle annihilation, pair production, particle decay etc. It is, of course, just a convention of language that we visualise particles, i.e. the constituents of our own universe as moving "forward" in time, thus establishing a causal link with the "past" and not with the "future". We are thus led to the result that in atomic-scale reactions, forward motion in time of all the reactants/products is a necessary condition for conservation of energy. What happens if energy is not conserved ? The universe contains but two things — energy and massive matter — and mass$\leftrightarrow$energy conversion is the only possible outcome of such reactions. However, according to our earlier result, such "non-conservation" of energy in atomic-scale reactions can take place if and only if one or more of the reactants/products do not move forward in time, i.e. move backward in time (since time is one-dimensional). We may ask the question from the opposite viewpoint : what happens in a reaction if one or more of the reactants/products travel backward

in time ? The answer is simple; energy will not be conserved, hence mass↔energy conversion must occur in such reactions. (Of course, the precise time interval between the "start" of the reaction and the mass↔energy conversion depends on the particular characteristics in the process involved; for example, positron-electron annihilation can take place either instantaneously or after positronium formation, depending on the collision geometry.)

We see therefore that from our time's arrow, we have obtained the result that in reactions where mass↔energy conversion takes place, at least one of the reactant/product particles must be travelling in time in a direction opposite to the others. When we recall that such reactions must involve both "particles" and "antiparticles", then from known properties of antiparticles, we conclude that an antiparticle can be obtained only under a combined operation of charge conjugation C, parity P and time reversal T (in any order) on a particle; similarly, CPT operating on an antiparticle gives us a particle. Since the *total number* of particles and antiparticles in any process involving particles and antiparticles is always even, the matrix element for the process remains invariant under the CPT operation. (This invariance is not quite self-evident, as we remember that the CPT theorem is supposed to be valid for *local* field theories only; our work seems to have broadened its scope.)

To conclude, we have shown that matter fields possess a non-local U(1) gauge transformation symmetry that necessitates invoking a non-local electromagnetic field, and that a non-thermodynamic arrow of time at the quantum (atomic) level is enforced upon us as a necessary condition for conservation of energy in matter-radiation interactions. We have also predicted that as a necessary condition for energy conservation in matter-radiation interactions, all particles in the universe subject to electromagnetic interactions should follow this one and the same arrow of time through the agency of the CMBR. We have also found a physical basis for Feynman's idea that antiparticles can be considered as particles moving backward in time, and we have also found that the scope of the CPT theorem can be extended to cover non-local gauge fields.

APPENDIX A

Elementary analysis shows that

$$\mathcal{R} = \int_0^\infty R_i(r) r^2\, dr \int_0^\infty R_f(r) r\, dr,$$

$$\mathcal{I}_1 = \sum_{n=0}^{\infty} \frac{4\pi a_n}{2n+1} \sum_{\nu=-n}^{n} \int Y^*_{l_f m_f}(\hat{r}) Y_{n\nu}(\hat{r})\, d\hat{r} \int Y^*_{n\nu}(\hat{r}') Y_{l_i m_i}(\hat{r}')\, d\hat{r}',$$

$$\text{and} \quad \mathcal{I}_2 = \sum_{n=0}^{\infty} \tfrac{16\pi^2 a_n}{3(2n+1)} \sum_{\nu=-n}^{n} \sum_{m=-1}^{1} \int Y^*_{l_f m_f}(\hat{r}) Y_{1m}(\hat{r}) Y_{n\nu}(\hat{r})\, d\hat{r} \times \int Y^*_{1m}(\hat{r}') Y^*_{n\nu}(\hat{r}') Y_{l_i m_i}(\hat{r}')\, d\hat{r}'.$$

The first angular integral gives $l_f = l_i$, while the second integral, which can be evaluated using the Gaunt formula, vanishes unless $n = l_f \pm 1$ and $n = l_i \pm 1$. Combining these two we get the selection rules $l_f = l_i,\ l_i \pm 2$.

APPENDIX B

The materials in this appendix are standard, and are included only for the sake of completeness. More details may be found in [9] (we have introduced some minor changes in notation). The Schrödinger equation for a two- electron atom in hyperspherical coordinates takes the form

$$\left[\frac{\hbar^2}{2m}\left(-\frac{\partial^2}{\partial R^2}+\frac{\Lambda^2}{R^2}\right)+V-E\right]\Psi(R,\alpha,\hat{r},\hat{r}')=0, \tag{D1}$$

where the "grand angular momentum operator"

$$\Lambda^2=-\frac{\partial^2}{\partial\alpha^2}-\frac{1}{4}+\frac{l(l+1)}{\cos^2\alpha}+\frac{l'(l'+1)}{\sin^2\alpha}. \tag{D2}$$

Also, $V=-\frac{Z}{r}-\frac{Z}{r'}+\frac{1}{|\mathbf{r}-\mathbf{r}'|}$, the total Coulomb interaction within the atom, and m is the electron mass. The channel eigenfunctions Φ's in (20) are obtained by configuration mixing in a basis of eigenfunctions of the grand angular momentum operator (D2). The eigenfunctions of this operator are

$$\phi_{nll'LM}=A(\cos\alpha)^l(\sin\alpha)^{l'}P_n^{(a,b)}(\cos 2\alpha)\mathcal{Y}^M_{Lll'}(\hat{r},\hat{r}'), \tag{D3}$$

where A is a normalisation constant, $a=l'+\frac{1}{2}$, $b=l+\frac{1}{2}$,

$$\mathcal{Y}^M_{Lll'}=\sum_{m=-l}^{l}\sum_{m'=-l'}^{l'}(ll'mm'|ll'LM)Y_{lm}(\hat{r})Y_{l'm'}(\hat{r}'),$$

$$P_n^{(a,b)}(x)=\tfrac{(a+1)_n}{n!}\,{}_2F_1\left(-n,a+b+n+1;a+1;\tfrac{1-x}{2}\right) \tag{D4}$$

is the Jacobi polynomial, and n is the radial correlation quantum number, the same as n_{rc} of Fano [9]. The eigenvalues of Λ^2 are given by $\lambda(\lambda+1)\hbar^2$, where $\lambda=2n+l+l'+\frac{3}{2}$. Thus for each channel μ in (20) we can write

$$\Phi_\mu=\sum_j C_{\mu j}\phi_j(\alpha,\hat{r},\hat{r}'),$$

where the running index $j\equiv(n,l,l',L,M)$ collectively.

In hyperspherical coordinates the total Coulomb potential V is given by -

$$V=\frac{e^2}{R}C(\alpha,\hat{r},\hat{r}'),\quad C=-\frac{Z}{\cos\alpha}-\frac{Z}{\sin\alpha}+\frac{1}{(1-\hat{r}.\hat{r}'\sin 2\alpha)^{1/2}}. \tag{D5}$$

The hyperradial functions $F_\mu(R)$ in (20) are given by the coupled equations

$$\left[\frac{\hbar^2}{2m}\left(-\frac{d^2}{dR^2}+\frac{\lambda(\lambda+1)}{R^2}\right)-E\right]F_\mu(R)+\frac{e^2}{R}\sum_{\mu'}V_{\mu\mu'}F_{\mu'}(R)\delta_{LL'}\delta_{MM'}=0. \tag{D6}$$

Restricting ourselves to a single configuration in (20) and substituting eqn.(16) in (19), we obtain eqn. (21) for $T_{\rm II}$.

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