

Classically Quantised Vacuum Electrodynamics, Part III: Vortical Photons as Maximally Entangled Pairs

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Abstract

Using a classical approach, we identify the electromagnetic fields of maximally entangled photon pairs. This construction is based on the excitation sector established in the companion article, in which displacement is the primitive quantised excitation of a single spatial volume, or voxel. Within that voxel, the electromagnetic field components satisfy the algebraic closure $\mathfrak{M}$. Its non-trivial one-axis solutions describe electromagnetic vortices travelling at the speed of light along their centrelines.

The nil-residue assumption introduced in Part I is applied as a global deterministic constraint. For an idealised recoilless emission of two equal-frequency, circularly polarised photons propagating in opposite directions with opposite gyrarities, their combined fields satisfy

$$\mathfrak{L}_{12}(t) \propto (\mathbf{B}_1 + \mathbf{B}_2) \times (\mathbf{E}_1 + \mathbf{E}_2) = 0.$$

Preservation of this nil residue correlates a later linear polarisation of one photon with the orthogonal polarisation of its partner. Application of Malus' law gives

$$E(\theta_1, \theta_2) = \cos(2(\theta_1 - \theta_2)),$$

the experimentally observed polarisation correlation characteristic of a maximally entangled photon pair. This correlation is obtained without representing either photon as an indeterminate Hilbert-space superposition; the construction is deterministic and explicitly nonlocal.

The action of a general retarding wave plate is also derived within the same photon parameterisation. For arbitrary incident polarisation, plate angle, retardance, running phase, and helicity, the transformed Maxwellian field agrees pointwise with the real polarisation trace obtained from Jones algebra.

1 Introduction

The companion article [1] (Part I) — whose results are cited with the prefix I — introduced the cyclic algebraic closure $\mathfrak{M}$ (Axiom I.3.2) and its Maxwellian instantiation on $(\mathbf{v}, \mathbf{B}, \mathbf{E})$,

$$\mathcal{M}[\mathbf{v}, \mathbf{B}, \mathbf{E}] := \left\{ \mathbf{v}, \mathbf{B}, \mathbf{E} \in \mathcal{S} \mid \mathbf{v} = \frac{\mathbf{B} \times \mathbf{E}}{\mathbf{B} \cdot \mathbf{B}}, \quad \mathbf{B} = \frac{\mathbf{E} \times \mathbf{v}}{\mathbf{v} \cdot \mathbf{v}}, \quad \mathbf{E} = \mathbf{v} \times \mathbf{B} \right\}, \quad (1)$$

over the complex vector space $\mathcal{S} \subset \mathbb{C}^3$, equipped with a non-Hermitian symmetric bilinear product. Within the vacuum-transverse sector of Assumption I.2.1, the closure forces the triple to

be pairwise orthogonal (Corollary I.3.3); under inverse Lie transport it reconstructs the source-free Maxwell curl equations (Proposition I.3.5). Part I further identified the one-axis travelling gyration as the basis for a photon (Assumption I.3.6), with gyricity $\mathfrak{g}$, directicity $\mathfrak{d}$, and dimensionless helicity $\mathfrak{h} = \mathfrak{g}\mathfrak{d}$. It also introduced the excitation sector, in which Maxwell's electric displacement is the primitive excitation agent of a quantised disturbance propagating through the vacuum (Hypothesis I.4.7).

The companion article [2] (Part II) — whose results are cited with the prefix II — recast Maxwell's 1865 equations for this excitation sector. It developed the corresponding displacement fields, currents, Beltrami fields and vector potentials, and derived the Einstein–Planck relation $\mathcal{E} = h\nu$ (Corollary II.6.1). Together, Parts I and II supply the kinematic, field-theoretic, and energetic foundations used here.

The present article develops a classical photon model from the one-axis solution of (1). It parameterises the complete polarisation state on a complex World-Sphere, together with gyricity, directicity, helicity, and optical parity. The nil-residue assumption introduced in Part I (Assumption I.2.4) is then specialised as a global deterministic constraint on an entangled photon pair. Applied to the calcium-cascade experiment, the model derives the observed two-channel correlation

$$E(\theta_A, \theta_B) = \cos(2(\theta_A - \theta_B))$$

without assigning an indeterminate superposition to either photon. Finally, the action of a general retarding wave plate is derived within the same photon parameterisation and shown to agree pointwise with Jones algebra for arbitrary plate angle, retardance, polarisation state, and gyricity.

The scope of the construction is the polarisation sector addressed by these results: the classical photon parameterisation, its two-photon entanglement correlation, and its transformation by ideal wave plates. Section 2 constructs the photon model; Section 3 develops the nil-residue account of entanglement; and Section 4 establishes its correspondence with Jones algebra. Appendix A proves the pointwise Jones–Maxwellian identity.

2 Classical Photon Model

Definition 2.1 (Photon Parameterisation). Below, we define the parameters to describe a photon in terms of the $\mathbf{v}$, $\mathbf{B}$, and $\mathbf{E}$ fields which are incorporated in the continuum-field realisation of (4) satisfying the Maxwellian closure $\mathfrak{M}$:

$\mathfrak{d} \in \{+1, -1\}$ defines directicity, the longitudinal propagation direction along the optical axis.

$\mathfrak{g} \in \{+1, -1\}$ defines gyricity, referenced to the laboratory frame.

$\mathfrak{h} = \mathfrak{d}\mathfrak{g}, \in \{+1, -1\}$ defines helicity.

$\vec{\mathfrak{s}}$ is the complex polarisation-state vector on the World-Sphere, defined by

$$\vec{\mathfrak{s}}(\alpha, \beta) = i \cos \alpha \sin \beta \hat{x} - i \sin \alpha \sin \beta \hat{y} + \cos \beta \hat{z}. \quad (2)$$

The image of (2) is therefore a real two-sphere embedded in the six-real-dimensional space $\mathbb{C}^3 \simeq \mathbb{R}^6$. More specifically, it lies in the real three-dimensional subspace spanned by $i\hat{x}$, $i\hat{y}$, and $\hat{z}$. We refer to this embedded two-sphere as the *complex World-Sphere* (CWS). The term is specific to the present construction and is not intended to identify

this surface with the conventional Poincaré or Bloch sphere. Here,
 $\beta \in \mathbb{R} \mid 0 \leq \beta \leq 2\pi$ defines the ellipticity of the polarisation; $\beta \in \{0, \pi\}$ correspond to a fully circularised polarisation states at the poles of CWS, whereas $\beta \in \{\pi/2, 3\pi/2\}$ correspond to fully linear polarised states on the CWS's equator. While,
 $\alpha \in \mathbb{R} \mid 0 \leq \alpha \leq 2\pi$ specifies the transverse orientation parameter.
 $\mathbf{p} = e^{i2\vartheta}$ is the parity operator, where ϑ represents the angle of the reflection axis relative to the reference frame.

With these definitions, we obtain a complete classical parameter set on the World-Sphere:

$$\mathcal{M} = \left\{ \vec{\mathbf{s}}(\alpha, \beta), \quad \mathbf{p}(\vartheta) = e^{i2\vartheta}, \quad \mathbf{d} \in \{+1, -1\}, \quad \mathbf{g} \in \{+1, -1\} \right\}. \quad (3)$$

This set completely tracks transverse polarisation orientation ($\vec{\mathbf{s}}$), optical reflection/rotation ($\mathbf{p}$), longitudinal propagation direction ($\mathbf{d}$), and intrinsic field chirality ($\mathbf{g}$) without needing complex Hilbert space state vectors or density matrices.

Remark 2.2 (Half-Angle Mapping to State Space). The full complex World-Sphere parametrization (α, β) directly maps to traditional Poincaré or Bloch parameter spaces via the half-angle substitution $\psi = \alpha/2 \pmod{\pi/2}$ and $\chi = \beta/2 \pmod{\pi/4}$. While conventional Poincaré representations quotient out the antipodal phase signs to simplify classical intensity measurements, maintaining the unquotiented coordinates (α, β) on the complex World-Sphere preserves the full $\mathbb{Z}_2$ geometric phase continuity required to describe interference and entanglement transformations.

The problem to be solved here is to find the transformation from the elementary transverse representation to a representation that describes a photon's electromagnetic field structure for any combination of parameters presented by $\mathcal{M}$, defined by (3), represented below as unit vectors:

$$\begin{bmatrix} \hat{\mathbf{v}} \\ \hat{\mathbf{B}} \\ \hat{\mathbf{E}} \end{bmatrix} = I \overbrace{\begin{bmatrix} 0 & 0 & 1 \\ C_\tau & -S_\tau & 0 \\ S_\tau & C_\tau & 0 \end{bmatrix}}^{R_0} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \hat{z} \end{bmatrix} \xrightarrow{\mathcal{M}} \text{diag} \begin{bmatrix} 1 \\ 1 \\ \mathbf{d} \end{bmatrix} \overbrace{\begin{bmatrix} 0 & 0 & \mathbf{d} \\ \mathcal{F} & -\mathcal{G} & 0 \\ \mathcal{G} & \mathcal{F} & 0 \end{bmatrix}}^{R_{\mathcal{M}}} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \mathbf{p}\hat{z} \end{bmatrix} \quad (4)$$

where $\mathcal{F}$ and $\mathcal{G}$ are complex functions of the parameters set $\mathcal{M}$, and I the unit matrix.

Furthermore, the closure $\mathfrak{M}$ dictates that the transversality condition of (4) is always preserved. The complex photon construction is obtained from an ordinary real rotation as a precursor. Under polarisation, the real gyration xy -plane of R_0 is first rotated in $\mathbb{R}^3$, using a three dimensional Tait-Bryan rotation

$$\begin{bmatrix} \vec{\mathbf{s}}_R \\ \hat{\mathbf{B}}' \\ \hat{\mathbf{E}}' \end{bmatrix} = \text{diag} \begin{bmatrix} 1 \\ \hat{\mathbf{B}} \\ \hat{\mathbf{E}} \end{bmatrix} \begin{pmatrix} C_\alpha S_\beta & -S_\alpha S_\beta & C_\beta \\ C_\tau C_\alpha C_\beta - S_\tau S_\alpha & -C_\tau S_\alpha C_\beta - S_\tau C_\alpha & -C_\tau S_\beta \\ S_\tau C_\alpha C_\beta + C_\tau S_\alpha & -S_\tau S_\alpha C_\beta + C_\tau C_\alpha & -S_\tau S_\beta \end{pmatrix} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \hat{z} \end{bmatrix}. \quad (5)$$

The first row,

$$\vec{\mathbf{s}}_R = C_\alpha S_\beta \hat{\mathbf{x}} - S_\alpha S_\beta \hat{\mathbf{y}} + C_\beta \hat{\mathbf{z}}, \quad (6)$$

is the real rotational precursor of the complex World-Sphere state. Placing its transverse compo-

nents in quadrature gives

$$\vec{\mathfrak{s}}_R \longmapsto \vec{\mathfrak{s}} = i C_\alpha S_\beta \hat{\mathbf{x}} - i S_\alpha S_\beta \hat{\mathbf{y}} + C_\beta \hat{\mathbf{z}},$$

which is precisely (2).

The same geometrical complexification is applied to the $\hat{\mathbf{B}}'$ - $\hat{\mathbf{E}}'$ gyration plane. The $\hat{\mathbf{z}}$ excursion of $\hat{\mathbf{B}}'$ is transferred into quadrature components along $\hat{x}$ and $\hat{y}$ in the ratio $\cos \alpha : \sin \alpha$. This gives

$$\begin{aligned} \hat{\mathbf{B}} = & (C_\tau C_\alpha C_\beta - i C_\tau C_\alpha S_\beta - S_\tau S_\alpha) \hat{\mathbf{x}} \\ & + (-C_\tau S_\alpha C_\beta - i C_\tau S_\alpha S_\beta - S_\tau C_\alpha) \hat{\mathbf{y}}, \end{aligned}$$

and hence

$$\begin{aligned} \mathcal{F} &:= e^{-i\beta} C_\tau C_\alpha - S_\tau S_\alpha, \\ \mathcal{G} &:= e^{i\beta} C_\tau S_\alpha + S_\tau C_\alpha. \end{aligned} \quad (7)$$

Substituting this results back into (4), we obtain classical description of a photon parameterised by $\mathcal{M}$ defined in (3)

$$\begin{bmatrix} \mathbf{v} \\ \mathbf{B} \\ \mathbf{E} \end{bmatrix} = \text{diag} \begin{bmatrix} c \\ B \\ \mathfrak{d}cB \end{bmatrix} \begin{bmatrix} 0 & 0 & \mathfrak{d} \\ e^{-i\beta} C_\tau C_\alpha - \mathfrak{g} S_\tau S_\alpha & -e^{i\beta} C_\tau S_\alpha - \mathfrak{g} S_\tau C_\alpha & 0 \\ e^{i\beta} C_\tau S_\alpha + \mathfrak{g} S_\tau C_\alpha & e^{-i\beta} C_\tau C_\alpha - \mathfrak{g} S_\tau S_\alpha & 0 \end{bmatrix} \begin{bmatrix} \hat{\mathbf{x}} \\ \hat{\mathbf{y}} \\ \mathfrak{p}\hat{\mathbf{z}} \end{bmatrix}. \quad (8)$$

Definition 2.3 (Photon spin). Within the present construction, the dimensionless photon spin vector is defined by

$$\vec{S} := \mathfrak{d} \mathfrak{g} \vec{\mathfrak{s}} = \mathfrak{d} \mathfrak{g} (i C_\alpha S_\beta \hat{\mathbf{x}} - i S_\alpha S_\beta \hat{\mathbf{y}} + C_\beta \hat{\mathbf{z}}). \quad (9)$$

Since $\mathfrak{d}^2 = \mathfrak{g}^2 = 1$ and $\vec{\mathfrak{s}}^\dagger \vec{\mathfrak{s}} = 1$ gives $\vec{S}^\dagger \vec{S} = 1$. Thus $\vec{S}$ has unit Hermitian norm for every polarisation state. The real projection of (9) is

$$\Re \vec{S} = \mathfrak{d} \mathfrak{g} \cos \beta \hat{\mathbf{z}}. \quad (10)$$

For a single photon the corresponding measurable longitudinal spin angular momentum is therefore

$$\mathbf{S}_{\text{SAM}} = \hbar \Re \vec{S} = \hbar \mathfrak{d} \mathfrak{g} \cos \beta \hat{\mathbf{z}}. \quad (11)$$

Remark 2.4 (Linear and circular polarisation). For linear polarisation, $\beta = \pi/2$, and

$$\Re \vec{S} = 0, \quad \mathbf{S}_{\text{SAM}} = 0.$$

Thus a linearly polarised photon has no measurable longitudinal spin angular momentum in the present construction, although the underlying complex spin vector retains unit Hermitian norm. At the circular pole, $\beta = 0$, one instead obtains

$$\mathbf{S}_{\text{SAM}} = \hbar \mathfrak{d} \mathfrak{g} \hat{\mathbf{z}},$$

so that the sign of the observable angular momentum is fixed jointly by the propagation direction and intrinsic handedness.

The mechanical transfer of optical spin angular momentum has direct experimental precedent. Poynting analysed the torque associated with changing circular polarisation by means of

a birefringent plate [3], and Beth subsequently measured the corresponding mechanical torque produced by changes in the angular momentum of polarised light [4].

Remark 2.5 (On the chirality factor $\mathfrak{g}$). If we map $\tau \mapsto \omega\tau$, then from a purely geometric and trigonometrical view chirality $\mathfrak{g} := \text{sign } \omega$. This changes neither the energy nor the momentum, both involving the Planck–Einstein relation $E = \hbar\omega$ [5, 6]. Historically introduced as a quantum hypothesis, but now derived in the excitation sector, see II.6.1 (Energy content of a quantised disturbance $\mathcal{W}$). Energy is proportional to $B^2 + E^2$ and is thus always positive; energy and linear-momentum calculations therefore remain unchanged, while the sign of the angular momentum is determined by $\mathfrak{g}$.

Remark 2.6 (On the parity factor $\mathfrak{p}$). In (8), $\hat{\mathbf{B}}$ and $\hat{\mathbf{E}}$ carry no $\hat{z}$ -component, so $\mathfrak{p}$ acts only on $\mathbf{v}$ at this stage; it enters $\mathbf{B}, \mathbf{E}$ only once an explicit reflection (e.g. $\hat{z} \rightarrow \hat{y}$) is applied via a further rotation of R . A detailed treatment of reflections is left out of scope here for brevity.

3 Photon Entanglement

Born’s statistical interpretation made probability fundamental to quantum mechanics [7]. The de Broglie–Bohm theory retained instead a deterministic dynamics in which definite configurations are guided by the wavefunction [8–10]. In the experimental domain considered here, the two descriptions reproduce the same outcome statistics. Their difference has therefore remained ontological: orthodox quantum mechanics assigns irreducible probabilities to measurement outcomes, whereas Bohmian mechanics attributes those outcomes to a definite, although nonlocal, dynamics. The present work asks whether a global constraint on Bohmian guidance identifies an optical regime in which the two accounts make different predictions.

Bell’s theorem converted part of the foundations debate into an experimental question [11], and the CHSH inequality gave it a form suited to laboratory tests [12]. Freedman and Clauser obtained the first experimental violation of a Bell inequality from the polarisation correlations of a calcium atomic cascade [13]. Aspect and co-workers developed two-channel measurements and later varied the analyser settings during photon flight [14, 15].

These experiments exploited the two-photon cascade of the excited calcium atom from the upper level $4p^2\ ^1S_0$ to the ground state $4s^2\ ^1S_0$ (both at $J = 0$) via the intermediate $J = 1$ state $4s4p\ ^1P_1$, emitting two photons in succession: ν_1 at 551.3 nm (green) and ν_2 at 422.7 nm (violet).

Angular-momentum conservation across this $J = 0 \rightarrow 1 \rightarrow 0$ sequence permits only two decay routes, and these fix, at emission, a definite pair of circular-polarisation values: either both photons right-handed or both left-handed. Taken on its own, this correlation is unremarkable and entirely local — the calcium atom is doing no more than Bell’s own Bertlmann, who decides locally, once, which colour socks to wear, so that finding one sock pink tells you the other is not pink [16]. The same holds for the frequency correlation: whichever photon is found green, the partner is violet, fixed by the cascade’s energy levels at emission, not by the act of measurement.

Entanglement proper is not this same-basis matching but its survival under a change of basis. If each photon is instead analysed by a two-channel linear polariser — in Aspect’s realisation, a polarising beam splitter (PBS) set at angle θ , giving simultaneous $+/-$ outcomes [15] — quantum mechanics predicts the experimentally confirmed correlation

$$E(\theta_A, \theta_B) = \cos(2(\theta_A - \theta_B)), \quad (12)$$

which no single value carried locally by each photon from the moment of emission can reproduce for every independent choice of θ_A, θ_B : this is the content of Bell's theorem [11]. The calcium atom, at the instant it emits ν_1 and ν_2 , cannot know which analyser angle will later be chosen independently at each wing; whatever fixes the outcome pair correctly for every such choice cannot therefore be a local property assigned once and carried away by each photon — on the deterministic hypothesis it requires a global hidden variable linking the pair as a whole, not two independently decided local ones. Determinism survives this argument intact; what is excluded is a *local* determinism of the Bertlmann-sock kind. We identify the global hidden variable with the operation of the nil-residue universe introduced in Part I. Assumption I.2.4 requires the universal ledger to vanish componentwise at every instant (Equation I.3) and requires variations in every partition of the closed universe to cancel (Equation I.5). Here we specialise that general assumption to a two-photon interaction.

Because this balance must hold at every instant and for every partition, its operation cannot in general be mediated by a local, subluminal signal between the two regions. The constraint applies to the universe as a whole and only afterward decomposes into local contributions. In the deterministic account pursued here, it supplies the global hidden variable permitted by Bell's theorem [11]: the outcome pair is fixed rather than created at measurement, while locality is not claimed. The assumption is to be judged here by the correlation obtained from its application to the photon fields.

We now examine an idealised counterpart of the calcium-cascade experiment of Freedman and Clauser [13], using Assumption I.2.4 (Nil-residue universe), the closure $\mathfrak{M}$, and the photon model developed in Section 2. The physical cascade is a three-body process involving two photons and the recoiling atom. Here we consider a recoilless event in which two equal-frequency photons, $\nu_1 = \nu_2 = \nu$, are emitted with opposite directivities and opposite gyricities. Their linear and angular momenta therefore cancel.

Theorem 3.1 (Classical entanglement by nil residue). *In a recoilless event, an atom emits two equal-frequency photons with opposite directivities and opposite gyricities. Their running phases are separated by $\pi/2$, as specified in (13). Define their joint action ledger using the non-Hermitian symmetric bilinear product of the closure $\mathfrak{M}$:*

$$\mathfrak{L}_{12}(t_0) \propto (\hat{\mathbf{B}}_1 + \hat{\mathbf{B}}_2) \times (\hat{\mathbf{E}}_1 + \hat{\mathbf{E}}_2) = 0.$$

If photon 1 is subsequently linearly polarised on an axis θ_1 , preservation of nil residue determines the orthogonal linear polarisation of photon 2, together with the gyricity-dependent phase reparameterisation required by (8). The resulting bilinear action ledger vanishes identically:

$$\mathfrak{L}_{12}(t) = 0.$$

Under Malus' law, a subsequent two-channel analysis of photon 2 on an axis θ_2 gives

$$P_{\parallel}(\theta_1, \theta_2) = \sin^2(\theta_1 - \theta_2), \quad P_{\perp}(\theta_1, \theta_2) = \cos^2(\theta_1 - \theta_2),$$

and hence

$$E(\theta_1, \theta_2) := P_{\perp} - P_{\parallel} = \cos(2(\theta_1 - \theta_2)).$$

Thus, within the classical Maxwellian photon model, the nil-residue constraint reproduces the polari-

sation correlation of the two-photon state.

Proof. We model the two photons, ν_1 and ν_2 , emitted by the idealised calcium cascade using the unit-vector representation of (8). At $t = t_0$, impose the initial conditions

$$\beta_1 = \beta_2 = \alpha_1 = \alpha_2 = 0, \quad \mathbf{d}_1 = -\mathbf{d}_2, \quad \mathbf{g}_1 = -\mathbf{g}_2,$$

so that

$$\mathbf{d}_1 \mathbf{g}_1 = \mathbf{d}_2 \mathbf{g}_2.$$

Let the phase difference between the two photons be $\pi/2$. Thus, writing

$$c_1 = \cos(\nu t), \quad s_1 = \sin(\nu t),$$

and

$$c_2 = \cos(\nu t + \pi/2), \quad s_2 = \sin(\nu t + \pi/2),$$

we obtain

$$\begin{bmatrix} \hat{\mathbf{v}}_1 \\ \hat{\mathbf{B}}_1 \\ \hat{\mathbf{E}}_1 \end{bmatrix} = \text{diag} \begin{bmatrix} 1 \\ 1 \\ \mathbf{d}_1 \end{bmatrix} \begin{bmatrix} 0 & 0 & \mathbf{d}_1 \\ c_1 & -\mathbf{g}_1 s_1 & 0 \\ \mathbf{g}_1 s_1 & c_1 & 0 \end{bmatrix} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \hat{z} \end{bmatrix} \quad \begin{bmatrix} \hat{\mathbf{v}}_2 \\ \hat{\mathbf{B}}_2 \\ \hat{\mathbf{E}}_2 \end{bmatrix} = \text{diag} \begin{bmatrix} 1 \\ 1 \\ \mathbf{d}_2 \end{bmatrix} \begin{bmatrix} 0 & 0 & \mathbf{d}_2 \\ c_2 & -\mathbf{g}_2 s_2 & 0 \\ \mathbf{g}_2 s_2 & c_2 & 0 \end{bmatrix} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \hat{z} \end{bmatrix} \quad (13)$$

Theorem I.36 (Derived Electromagnetic Quantities), together with Equation I.36, gives

$$h = \epsilon_0 \phi \psi.$$

In vector form, the corresponding action is

$$\vec{h} = \epsilon_0 \vec{\phi} \times \vec{\psi}.$$

Because the magnetic and electric fields have the same directions as their respective fluxes, we define the nil-residue action ledger of the two photons at creation by

$$\mathfrak{L}_{12}(t) \propto (\hat{\mathbf{B}}_1 + \hat{\mathbf{B}}_2) \times (\hat{\mathbf{E}}_1 + \hat{\mathbf{E}}_2), \quad \mathfrak{L}_{12}(t_0) = 0. \quad (14)$$

At $t = t_1$, let photon 1 acquire a linear polarisation state, so that $\beta_1 = \pi/2$ while $\alpha_1 = 0$ remains unchanged. Equation (8) then gives

$$\begin{bmatrix} \hat{\mathbf{v}}_1 \\ \hat{\mathbf{B}}_1 \\ \hat{\mathbf{E}}_1 \end{bmatrix} = \text{diag} \begin{bmatrix} 1 \\ 1 \\ \mathbf{d}_1 \end{bmatrix} \begin{bmatrix} 0 & 0 & \mathbf{d}_1 \\ -ic_1 & -\mathbf{g}_1 s_1 & 0 \\ \mathbf{g}_1 s_1 & -ic_1 & 0 \end{bmatrix} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \hat{z} \end{bmatrix} \quad (15)$$

This local change would give $\mathfrak{L}_{12}(t_1) \neq 0$ if the state of photon 2 remained unchanged. Such a residue is prohibited by Assumption I.2.4. Preservation of the joint ledger therefore changes the polarisation of photon 2 at a distance, placing it in the orthogonal linear state

$$\beta_2 = \frac{\pi}{2}, \quad \alpha_2 = \frac{\pi}{2},$$

together with the gyricity-dependent phase reparameterisation required by (8). This gives

$$\begin{bmatrix} \hat{\mathbf{v}}_2 \\ \hat{\mathbf{B}}_2 \\ \hat{\mathbf{E}}_2 \end{bmatrix} = \text{diag} \begin{bmatrix} 1 \\ 1 \\ \mathfrak{d}_2 \end{bmatrix} \begin{bmatrix} 0 & 0 & \mathfrak{d}_2 \\ c_2 & -i\mathfrak{g}_2 s_2 & 0 \\ i\mathfrak{g}_2 s_2 & c_2 & 0 \end{bmatrix} \begin{bmatrix} \hat{x} \\ \hat{y} \\ \hat{z} \end{bmatrix} \quad (16)$$

Direct evaluation using the symmetric bilinear product restores

$$\mathfrak{L}_{12}(t_1 + t_{\mathfrak{g}}) = 0, \quad t_{\mathfrak{g}} = 0.$$

This is an in-flight preservation of the joint photon state, occurring before either photon is measured and converted into an electrical detector impulse. No energy is transferred between the photons during this preservation.

Let the orthogonal polarisation states in (15) and (16) be defined relative to the optical axis θ_1 . At a later time $t_2 > t_1$, photon 2 is analysed on an optical axis θ_2 . Its polarisation is offset from that analyser axis by

$$\Delta\theta = (\theta_1 - \theta_2) + \frac{\pi}{2}.$$

Malus' law therefore gives

$$P_{\parallel}(\theta_1, \theta_2) = \cos^2 \left[(\theta_1 - \theta_2) + \frac{\pi}{2} \right] = \sin^2(\theta_1 - \theta_2),$$

whereas the orthogonal-channel probability is

$$P_{\perp}(\theta_1, \theta_2) = \cos^2(\theta_1 - \theta_2).$$

Defining the correlation coefficient as the normalised difference between the orthogonal and parallel coincidence rates gives

$$\begin{aligned} E(\theta_1, \theta_2) &= P_{\perp}(\theta_1, \theta_2) - P_{\parallel}(\theta_1, \theta_2) \\ &= \cos^2(\theta_1 - \theta_2) - \sin^2(\theta_1 - \theta_2) \\ &= \cos(2(\theta_1 - \theta_2)). \end{aligned}$$

This is the characteristic polarisation correlation of a maximally entangled photon pair. □

The correlation obtained in Theorem 3.1 is the same correlation observed in the calcium-cascade experiments and their subsequent two-channel refinements cited above. The theorem is therefore empirically corroborated. Because this result follows from nil-residue preservation within the classical Maxwellian photon model, that agreement provides direct empirical support for Assumption I.2.4 in the two-photon polarisation domain.

Our ignorance does not justify declaring that each photon physically occupies both the left- and right-circular states. The present construction assigns definite parameters to each photon. What remains unknown is the exact phase at photon creation relative to our time reference, and consequently the phase of photon 1 at the polarisation event $t = t_1$. This unresolved phase prevents our coarse-grained macroscopic description from predicting an individual detector outcome, but it does not make the underlying event indeterministic.

Quantum mechanics converts this lack of knowledge into an ontological postulate: the unknown alternatives are represented as a physical superposition, and the resulting probabilities

are then attributed to Nature itself. The present construction rejects that inversion. The ignorance belongs to the observer, whose time resolution does not reveal the determining phase; it is not a physical property of the photon. To elevate that ignorance into fundamental indeterminacy is to impose a limitation of the human description upon Nature and then treat that imposed limitation as an explanation.

Hicks stated the methodological distinction plainly:

“A statement of a law is either a confession of ignorance or a mnemonic convenience.”
— W. M. Hicks [17]

In the present account, probabilistic superposition is a confession of unresolved phase information. It is not required as a description of the underlying physical state. Schrödinger expressed the broader objection forcefully:

“Let me say at the outset, that in this discourse, I am opposing not a few special statements of quantum mechanics held today, I am opposing as it were the whole of it, I am opposing its basic views that have been shaped 25 years ago, when Max Born put forward his probability interpretation, which was accepted by almost everybody. It has been worked out in great detail to form a scheme of admirable logical consistency that has been inculcated ever since to every young student of theoretical physics.”
— E. Schrödinger [18]

Bell, whose theorem establishes the impossibility of a local account of these correlations but not the impossibility of determinism, stated the methodological point directly:

“To show us that vagueness, subjectivity, and indeterminism, are not forced on us by experimental facts, but by deliberate theoretical choice?”
— J. S. Bell [19]

Theorem 3.1 supplies such an alternative in the two-photon polarisation domain. It accepts the nonlocality required by Bell’s theorem while rejecting the claim that the observed correlation requires either subjective measurement collapse or fundamental indeterminism.

Having recovered the experimentally observed quantum-mechanical polarisation correlation from the classical photon model, we now test the model against the standard optical description of polarisation transformations. The following section derives the action of a general retarding wave plate on the parameterised Maxwellian photon and compares the resulting field trajectory directly with Jones algebra.

4 Wave plate action

Using the classical parameterised model, and assuming the magnetic slow axis is aligned with the $\hat{y}$ at zero calibration, the wave-plate induced phasing to (3) is modeled by

$$\alpha \mapsto \alpha_\theta := \alpha - \theta_{\mathcal{M}}. \quad (17)$$

Here $\theta_{\mathcal{M}}$ denotes the signed wave-plate offset in the row-vector convention of the present construction.

Remark 4.1 ($\theta_{\mathcal{M}}$ compared with Jones column-vector θ_J). The Jones column-vector convention used below employs the oppositely signed rotation parameter,

$$\theta_J = -\theta_{\mathcal{M}}. \quad (18)$$

This sign relation reflects the respective row- and column-vector rotation conventions and does not represent a different physical orientation of the wave plate.

The slow-axis delay δ is accounted for in the slow-axis temporal component of (8), which gives the solution matrix

$$R_{\mathcal{M}}(\delta) = \begin{bmatrix} 0 & 0 & \mathbf{d} \\ e^{-i\beta} C_{\tau} C_{\alpha} - \mathfrak{h} S_{\tau} S_{\alpha} & -e^{i\beta} C[\tau - \delta] S_{\alpha} - \mathfrak{h} S[\tau - \delta] C_{\alpha} & 0 \\ -e^{i\beta} C[\tau - \delta] S_{\alpha} - \mathfrak{h} S[\tau - \delta] C_{\alpha} & e^{-i\beta} C_{\tau} C_{\alpha} - \mathfrak{h} S_{\tau} S_{\alpha} & 0 \end{bmatrix} \quad (19)$$

where $S[\tau - \delta]$ reads $\sin(\tau - \delta)$, with α evaluated throughout at the shifted value α_{θ} of (17). The output photon referred to the laboratory frame requires a further rotation $\theta_{\mathcal{M}}$. Using rotation matrix R_z with $\psi \mapsto \theta_{\mathcal{M}}$ provides the solution

$$\begin{bmatrix} \mathbf{v} \\ \mathbf{B} \\ \mathbf{E} \end{bmatrix} = \text{diag} \begin{bmatrix} c \\ B \\ \mathbf{d}cB \end{bmatrix} \begin{bmatrix} R_{\mathcal{M}}(\delta) \\ \\ \end{bmatrix} \begin{bmatrix} R_z(\theta_{\mathcal{M}}) \\ \\ \end{bmatrix} \begin{bmatrix} \hat{\mathbf{x}} \\ \hat{\mathbf{y}} \\ \mathbf{p}\hat{\mathbf{z}} \end{bmatrix}. \quad (20)$$

Proposition 4.2 (Exact correspondence with Jones algebra). *For comparison with the magnetic-vector construction $\hat{\mathbf{B}}(\tau, \alpha, \beta, \theta_{\mathcal{M}}, \delta, \mathfrak{g})$, we use the conventional Jones representation of an elliptically polarised state [20]. Let*

$$\mathfrak{g} \in \{-1, +1\}$$

denote the handedness. The incident Jones vector is defined as

$$\mathbf{J}_{\text{in}}(\alpha, \beta, \mathfrak{g}) = \frac{1}{\sqrt{1 + \cos^2 \beta}} \begin{pmatrix} \sin \alpha + i\mathfrak{g} \cos \alpha \cos \beta \\ \cos \alpha - i\mathfrak{g} \sin \alpha \cos \beta \end{pmatrix}. \quad (21)$$

The normalisation has been chosen such that

$$\mathbf{J}_{\text{in}}^{\dagger} \mathbf{J}_{\text{in}} = 1.$$

A rotation through an angle θ_J is represented by

$$\mathbf{R}(\theta_J) = \begin{pmatrix} \cos \theta_J & -\sin \theta_J \\ \sin \theta_J & \cos \theta_J \end{pmatrix}, \quad (22)$$

and an ideal retarder having retardance δ , with its principal axes aligned with the coordinate axes, is represented by

$$\mathbf{W}(\delta) = \begin{pmatrix} 1 & 0 \\ 0 & e^{i\delta} \end{pmatrix}. \quad (23)$$

Thus, for a retarder whose principal axis is rotated through θ_J , the output Jones vector is

$$\mathbf{J}_{\text{out}} = \mathbf{R}(\theta_J) \mathbf{W}(\delta) \mathbf{R}(-\theta_J) \mathbf{J}_{\text{in}}. \quad (24)$$

Writing, as shorthand for the θ_J already introduced in (22), $C_\theta := \cos \theta_J$ and $S_\theta := \sin \theta_J$, the rotated retarder matrix may equivalently be written explicitly as

$$\mathbf{W}(\theta_J, \delta) = \begin{pmatrix} C_\theta^2 + e^{i\delta} S_\theta^2 & C_\theta S_\theta (1 - e^{i\delta}) \\ C_\theta S_\theta (1 - e^{i\delta}) & S_\theta^2 + e^{i\delta} C_\theta^2 \end{pmatrix}, \quad (25)$$

so that

$$\mathbf{J}_{\text{out}} = \begin{pmatrix} J_x \\ J_y \end{pmatrix}. \quad (26)$$

To recover the real Cartesian ellipse corresponding to the complex Jones vector, the common phase parameter τ is introduced only after the Jones transformation. The real parametrisation used here is

$$\mathbf{E}_J(\tau, \alpha, \beta, \theta_J, \delta, \mathbf{g}) = \begin{pmatrix} x_J(\tau) \\ y_J(\tau) \end{pmatrix} = \Re \{ \mathbf{J}_{\text{out}} e^{-i(\tau + \mathbf{g}\pi/2)} \}, \quad (27)$$

or, componentwise,

$$x_J(\tau) = \Re [J_x e^{-i(\tau + \mathbf{g}\pi/2)}], \quad (28)$$

$$y_J(\tau) = \Re [J_y e^{-i(\tau + \mathbf{g}\pi/2)}]. \quad (29)$$

The additional phase offset

$$+ \mathbf{g} \frac{\pi}{2} \quad (30)$$

in Eq. (27) has no effect on the polarisation state represented by the Jones vector. Conventional Jones polarisation is insensitive to a common global phase: only relative phases between its components affect the polarisation ellipse. Consequently, multiplication of the complete Jones vector by a common phase factor does not change its Jones polarisation state.

The offset $\mathbf{g}\pi/2$ is introduced here for a different purpose. The Jones construction is to be compared point for point with the explicitly phase-parametrised magnetic vector

$$\hat{\mathbf{B}}(\tau, \alpha, \beta, \theta_{\mathcal{M}}, \delta, \mathbf{g}),$$

extracted from (20). The otherwise immaterial global Jones phase must therefore be fixed when the complex Jones vector is converted into the real, τ -dependent ellipse. The choice $\tau + \mathbf{g}\pi/2$ aligns the temporal parametrisation of the Jones ellipse with that of $\hat{\mathbf{B}}$, allowing the two constructions to be compared at the same value of τ . It changes neither the geometric Jones ellipse nor any Jones-polarisation observable; it fixes only the phase origin used for the point-by-point comparison.

The agreement is algebraic rather than numerical. Fixing the free amplitude B in (20) to $B = 1/N_\beta$ (Appendix A), so that $\hat{\mathbf{B}}_{xy}$ carries the same normalisation as the unit Jones vector (21), the Jones trace and the real transverse magnetic-vector trace satisfy the identity

$$\mathbf{E}_J(\tau, \alpha, \beta, \theta_J, \delta, \mathbf{g}) = \Re \hat{\mathbf{B}}_{xy}(\tau, \alpha, \beta, \theta_{\mathcal{M}}, \delta, \mathbf{g}) \quad (31)$$

identically, for all

$$\tau, \alpha, \beta, \theta_{\mathcal{M}}, \delta \in \mathbb{R}, \quad \mathbf{g} \in \{-1, +1\}.$$

The agreement is thus pointwise in the running phase τ , not merely an equality of the resulting polarisation ellipses: the two constructions trace the same curve in the same sense at every instant, for arbitrary plate angle and retardance, and for both circular senses.

Proof. Appendix A. □

5 Conclusion

This article has developed a classical model of the photon from the one-axis travelling gyration admitted by the Maxwellian closure $\mathfrak{M}$. The resulting parameterisation describes the complete polarisation state on the complex World-Sphere, together with directicity $\mathfrak{d}$, gyricity $\mathfrak{g}$, helicity $\mathfrak{h} = \mathfrak{d}\mathfrak{g}$, and optical parity $\mathfrak{p}$. These quantities give the photon a definite electromagnetic field structure at every instant while preserving the transversality and cyclic relations imposed by the closure.

For a pair of oppositely directed photons with opposite gyricities, the linear and spin angular momenta cancel. Assumption I.2.4, introduced in Part I, strengthens this initial balance into a universe-wide constraint requiring the total bilinear action ledger to remain zero. Applied to the idealised two-photon cascade, preservation of that ledger correlates the polarisation acquired by one photon with the orthogonal polarisation of its partner. Together with Malus' law, this gives

$$E(\theta_1, \theta_2) = \cos(2(\theta_1 - \theta_2)),$$

the characteristic polarisation correlation of a maximally entangled photon pair. Its agreement with the calcium-cascade experiments empirically corroborates Theorem 3.1 and provides direct empirical support for the nil-residue assumption in the two-photon polarisation domain.

The resulting account is deterministic but explicitly nonlocal. The correlated outcome belongs to the photon pair under the global constraint and cannot be decomposed into two independently carried local instructions. The phase that determines an individual event remains unresolved by the macroscopic observer, but it is not therefore absent from the physical process. Within the domain examined here, probability measures this unresolved information; it need not be imposed upon Nature as an indeterminate physical superposition.

The same photon model also reproduces the established action of an ideal retarding wave plate. For arbitrary incident polarisation, plate angle, retardance, running phase, and gyricity, the transformed Maxwellian magnetic vector agrees pointwise with the real trace obtained from Jones algebra. The agreement therefore extends beyond the geometry of the resulting polarisation ellipse: both constructions describe the same trajectory, in the same sense, at every value of the running phase.

This result completes the sequence developed across the three articles. Part I introduced the cyclic algebraic closure $\mathfrak{M}$, recovered the source-free Maxwell equations from its transport, and identified the quantised gyration of the vacuum. Part II developed the excitation sector of that vacuum and obtained the Einstein–Planck relation $\mathcal{E} = h\nu$. The present article has carried those foundations to the photon itself, its complete polarisation parameterisation, its transformation by optical elements, and the experimentally observed correlation of maximally entangled pairs.

Taken together, the trilogy provides a continuous line from algebraic closure, through Maxwellian field dynamics and quantised excitation, to photon optics and entanglement. Phenomena usually assigned exclusively to Hilbert-space superposition are thereby recovered, within the scope established here, from definite classical electromagnetic configurations governed by a global nil-residue constraint.

A Exact Jones–Maxwellian wave-plate identity

We prove the pointwise identity (31). Define the componentwise real-trace operator

$$\mathcal{T}_{\mathfrak{g}}[\mathbf{z}](\tau) := \Re \left[\mathbf{z} e^{-i(\tau + \mathfrak{g}\pi/2)} \right]. \quad (32)$$

Since $\mathfrak{g} \in \{-1, +1\}$,

$$e^{-i\mathfrak{g}\pi/2} = -i\mathfrak{g}.$$

Writing

$$q := \cos \beta, \quad N_{\beta} := \sqrt{1 + q^2},$$

application of (32) to (21) gives

$$\mathcal{T}_{\mathfrak{g}}[\mathbf{J}_{\text{in}}](\tau) = \frac{1}{N_{\beta}} \begin{pmatrix} q \cos \tau \cos \alpha - \mathfrak{g} \sin \tau \sin \alpha \\ -q \cos \tau \sin \alpha - \mathfrak{g} \sin \tau \cos \alpha \end{pmatrix}. \quad (33)$$

Two elementary identities are sufficient to propagate this result through the wave plate. For any real matrix A ,

$$\mathcal{T}_{\mathfrak{g}}[A\mathbf{z}](\tau) = A \mathcal{T}_{\mathfrak{g}}[\mathbf{z}](\tau), \quad (34)$$

because taking the real part commutes with multiplication by a real matrix. Furthermore, for any complex scalar z ,

$$\mathcal{T}_{\mathfrak{g}}[e^{i\delta}z](\tau) = \mathcal{T}_{\mathfrak{g}}[z](\tau - \delta). \quad (35)$$

Thus the Jones retardance factor $e^{i\delta}$ is precisely a translation of the running temporal phase by δ .

The Jones rotation parameter and the MvBE wave-plate parameter have opposite signs according to (18),

$$\theta_J = -\theta_{\mathcal{M}}.$$

With the rotation convention of (22), direct multiplication gives the general identity

$$R(\varphi)\mathbf{J}_{\text{in}}(\alpha, \beta, \mathfrak{g}) = \mathbf{J}_{\text{in}}(\alpha - \varphi, \beta, \mathfrak{g}). \quad (36)$$

Hence

$$\begin{aligned} R(-\theta_J)\mathbf{J}_{\text{in}}(\alpha, \beta, \mathfrak{g}) &= \mathbf{J}_{\text{in}}(\alpha + \theta_J, \beta, \mathfrak{g}) \\ &= \mathbf{J}_{\text{in}}(\alpha - \theta_{\mathcal{M}}, \beta, \mathfrak{g}). \end{aligned} \quad (37)$$

Defining

$$\alpha_{\theta} := \alpha - \theta_{\mathcal{M}},$$

the Jones transformation into the wave-plate frame therefore uses exactly the same angular parameter as the MvBE construction.

Equations (33)–(36) then give

$$\begin{aligned}\mathbf{E}_J &= \mathcal{T}_{\mathbf{g}} [R(\theta_J)W(\delta)R(-\theta_J)\mathbf{J}_{\text{in}}](\tau) \\ &= \frac{1}{N_\beta} R(\theta_J) \begin{pmatrix} q \cos \tau \cos \alpha_\theta - \mathbf{g} \sin \tau \sin \alpha_\theta \\ -q \cos(\tau - \delta) \sin \alpha_\theta - \mathbf{g} \sin(\tau - \delta) \cos \alpha_\theta \end{pmatrix}.\end{aligned}\quad (38)$$

Now take the real, transverse part of the magnetic-vector row of $R_{\mathcal{M}}(\delta)$ in (19), with $\alpha \mapsto \alpha_\theta$. Since $\Re(e^{\pm i\beta}) = \cos \beta = q$, this gives directly

$$\Re \mathbf{B}_{xy} = \begin{pmatrix} q \cos \tau \cos \alpha_\theta - \mathbf{g} \sin \tau \sin \alpha_\theta \\ -q \cos(\tau - \delta) \sin \alpha_\theta - \mathbf{g} \sin(\tau - \delta) \cos \alpha_\theta \end{pmatrix}, \quad (39)$$

i.e. *exactly* the bracketed vector appearing in (38), with no factor of N_β at all.

This raw row is not yet the quantity compared in (31). It carries whatever overall amplitude the free constant B in the $\text{diag}(c, B, \mathbf{d}cB)$ prefactor of (20) happens to have, whereas the Jones vector (21) is fixed to unit norm by construction. That constant plays no role anywhere else in the construction, so we are free to fix it, once and for all, to

$$B := \frac{1}{N_\beta}, \quad (40)$$

matching the Jones normalisation exactly, and write $\hat{\mathbf{B}}_{xy} := \mathbf{B}_{xy}/N_\beta$ for the field row so normalised. This is a fixed choice of units, depending on β alone, made once at $\tau = 0$ -independent scale; it should not be confused with the separate, genuinely τ -dependent question, raised as an open item in Section ??, of whether the instantaneous Hermitian modulus $|\mathbf{B}(\tau)|^2$ itself stays constant under the wave plate — it does not, in general, and that fact is unrelated to, and unaffected by, the fixed overall normalisation introduced here.

With this choice, (39) gives precisely

$$\Re \hat{\mathbf{B}}_{xy} = \frac{1}{N_\beta} \begin{pmatrix} q \cos \tau \cos \alpha_\theta - \mathbf{g} \sin \tau \sin \alpha_\theta \\ -q \cos(\tau - \delta) \sin \alpha_\theta - \mathbf{g} \sin(\tau - \delta) \cos \alpha_\theta \end{pmatrix}. \quad (41)$$

Returning from the wave-plate frame to the laboratory frame by $R_z(\theta_{\mathcal{M}})$ in the row-vector convention is equivalent, after transposition, to multiplication by $R(\theta_J)$ in the Jones column-vector convention. Hence (41) becomes precisely (38), and therefore

$$\mathbf{E}_J(\tau, \alpha, \beta, \theta_J, \delta, \mathbf{g}) = \Re \hat{\mathbf{B}}_{xy}(\tau, \alpha, \beta, \theta_{\mathcal{M}}, \delta, \mathbf{g}).$$

The equality holds identically for arbitrary real $\tau, \alpha, \beta, \theta_{\mathcal{M}}, \delta$ and for both $\mathbf{g} = \pm 1$.

References

- [1] A. L. Vrba. *Classically Quantised Vacuum Electrodynamics: Maxwellian Vortices and an Emergent EM-Mass Gap*. 2026.

- [2] A. L. Vrba. *Classically Quantised Vacuum Electrodynamics: Maxwell's 1865 Monopole Recast and a Classical Derivation of the Einstein–Planck Relation*. 2026.
- [3] J. H. Poynting. “The Wave Motion of a Revolving Shaft, and a Suggestion as to the Angular Momentum in a Beam of Circularly Polarised Light”. In: *Proceedings of the Royal Society of London A* 82 (1909), pp. 560–567.
- [4] R. A. Beth. “Mechanical Detection and Measurement of the Angular Momentum of Light”. In: *Physical Review* 50 (1936), pp. 115–125.
- [5] M. Planck. “Ueber das Gesetz der Energieverteilung im Normalspectrum”. In: *Annalen der Physik* 309 (1901), pp. 553–563.
- [6] A. Einstein. “Über einen die Erzeugung und Verwandlung des Lichtes betreffenden heuristischen Gesichtspunkt”. In: *Annalen der Physik* 322 (1905), pp. 132–148.
- [7] M. Born. “Zur Quantenmechanik der Stoßvorgänge”. In: *Zeitschrift für Physik* 37 (1926), pp. 863–867.
- [8] L. de Broglie. “La Nouvelle Dynamique des Quanta”. In: *Électrons et Photons: Rapports et Discussions du Cinquième Conseil de Physique Tenu à Bruxelles du 24 au 29 Octobre 1927*. Paris: Gauthier-Villars, 1928, pp. 105–132.
- [9] D. Bohm. “A Suggested Interpretation of the Quantum Theory in Terms of “Hidden” Variables. I”. In: *Physical Review* 85 (1952), pp. 166–179.
- [10] D. Bohm. “A Suggested Interpretation of the Quantum Theory in Terms of “Hidden” Variables. II”. In: *Physical Review* 85 (1952), pp. 180–193.
- [11] J. S. Bell. “On the Einstein Podolsky Rosen Paradox”. In: *Physics Physique Fizika* 1 (1964), pp. 195–200.
- [12] J. F. Clauser et al. “Proposed Experiment to Test Local Hidden-Variable Theories”. In: *Physical Review Letters* 23 (1969), pp. 880–884.
- [13] S. J. Freedman and J. F. Clauser. “Experimental Test of Local Hidden-Variable Theories”. In: *Physical Review Letters* 28 (1972), pp. 938–941.
- [14] A. Aspect, P. Grangier, and G. Roger. “Experimental Realization of Einstein–Podolsky–Rosen–Bohm Gedankenexperiment: A New Violation of Bell’s Inequalities”. In: *Physical Review Letters* 49 (1982), pp. 91–94.
- [15] A. Aspect, J. Dalibard, and G. Roger. “Experimental Test of Bell’s Inequalities Using Time-Varying Analyzers”. In: *Physical Review Letters* 49 (1982), pp. 1804–1807.
- [16] J. S. Bell. “Bertlmann’s Socks and the Nature of Reality”. In: *Journal de Physique Colloques* 42.C2 (1981), pp. C2–41–C2–61.
- [17] W. M. Hicks. “Mathematics and Physics: Opening Address by Prof. W. M. Hicks”. In: *Nature* 52 (1895), pp. 472–477.
- [18] Erwin Schrödinger. *The Interpretation of Quantum Mechanics: Dublin Seminars (1949-1955 And Other Unpublished Essays)*. Ed. by Michel Bitbol. Ox-Bow Press, Connecticut, USA, 1995.
- [19] John S. Bell. “On the Impossible Pilot Wave”. In: *Foundations of Physics* 12.10 (1982), pp. 989–999.

- [20] R. C. Jones. “A New Calculus for the Treatment of Optical Systems. I. Description and Discussion of the Calculus”. In: *Journal of the Optical Society of America* 31 (1941), pp. 488–493.