The κ-Field Origin of $|\psi|^2$: Rate Law, Quadratic Measure, and Uniqueness

$\lambda_A = \kappa|\psi_A|^2$ from κ-Field Dynamics in the VERSF Framework

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For General Readers: What This Paper Is About

The question

Quantum mechanics tells us that the probability of measuring outcome A is $|\psi_A|^2$ — the squared magnitude of the complex amplitude. This is the Born rule, and it has been confirmed by every experiment ever performed. But why squared? Why not $|\psi_A|$ (the magnitude itself), or $|\psi_A|^4$ (the square of the square), or some other function? And where does this number go — how does it get from the quantum state into a physical detector?

This paper answers both questions from the same direction: the physics of how quantum amplitudes drive physical fields, and how physical detectors respond to those fields.

The answer in plain language

Every real detector operating without an external phase reference — a photodetector, a Geiger counter, a superconducting nanowire — responds to *power*, not to *amplitude*. Power is always a squared quantity. Sound intensity is the square of pressure amplitude. Light intensity is the square of electric field amplitude. This is not a quantum mechanical fact; it is a feature of how physical sensors work when they couple incoherently to a field bath with no reference signal.

Quantum amplitudes drive physical fields linearly — each possible outcome (each "branch") contributes to the mediating field in proportion to its amplitude ψ_A, not its amplitude squared. But when a detector measures that field, it measures power — the squared signal. The squaring happens at the detector, not in the quantum state.

After decoherence, the different branches cannot interfere. This means the cross-terms (branch A times branch B) vanish when we compute the power. What survives is a sum of independent squared contributions: $\Sigma_A |\psi_A|^2 \times$ (per-branch factor). The $|\psi_A|^2$ is therefore not a primitive of quantum mechanics but a consequence of three things acting together: linear sourcing, quadratic detector response, and decoherence eliminating interference.

This is also the *unique* form that survives. No other function of ψ_A — not $|\psi_A|$, not $|\psi_A|^4$, not any other exponent — is consistent with all four physical constraints: locality, analyticity, U(1) phase invariance, and independent branch contributions. The squaring is forced.

How this fits the larger programme

This paper is the companion to Taylor (2025a), which showed that if detector rates scale as $|\psi_A|^2$, the Born rule follows from first-passage competition between branches. That paper took the rate law as given; this paper derives it — establishing why rates must scale as $|\psi_A|^2$ from κ-field physics, and proving why the exponent is 2 and nothing else. Together, the two papers trace a complete physical chain from quantum amplitudes to measurement probabilities.

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Abstract

This paper addresses two related questions: why does the branch-dependent detector rate satisfy $\lambda_A = \kappa|\psi_A|^2$, and why is the exponent precisely 2 rather than any other value? Both are answered from the same physical argument, without invoking the Born rule as an axiom. The derivation proceeds in four steps. (i) The commitment current J decomposes in the branch basis with ψ_A entering as state-dependent amplitude weights at the operator level — the linearity is structurally forced by the linearity of quantum states and field equations. (ii) Physical detectors respond to field power, not field amplitude: every real detector measures energy, intensity, or

flux — all second-order in the field. The FGR transition rate in the Wigner–Weisskopf regime is proportional to the spectral density $S\_X(\omega\_cm)$, a second-order correlation function bilinear in the source. This is the step where the exponent shifts from 1 to 2. (iii) Decoherence diagonalises the noise power correlator — branches that have left distinguishable environmental records cannot coherently drive correlated κ-field fluctuations — giving $S\_X \propto \Sigma\_A |\psi\_A|^2 S\_A$. (iv) Branch-blindness equalises per-branch prefactors to within $N\_bath^{-1} \sim 10^{-23}$ corrections, yielding $\lambda\_A = \kappa|\psi\_A|^2$ with a single apparatus-dependent constant. Uniqueness is established at leading perturbative order by four constraints — locality, analyticity, U(1) covariance, and additivity — which admit only $|\psi\_A|^2$ from the polynomial expansion of any local analytic rate functional. The paper establishes a physical realisation result: the κ-field transmits the $|\psi\_A|^2$ structure from the quantum state to detector transition rates; it does not generate that structure from more primitive principles. The derivation assumes decoherence as established and constitutes a partial reduction: given decoherence, the Born rule follows from κ-field dynamics. A falsifiable prediction follows: detectors requiring $k = 2$ independent threshold events yield $P_1 = \lambda_1^2(\lambda_1 + 3\lambda_2)/(\lambda_1 + \lambda_2)^3 = 20/27 \approx 0.741$ for $\lambda_1 = 2\lambda_2$, deviating from the Born rule value of $2/3 \approx 0.667$ by approximately 0.074 (7.4 percentage points).

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1. Scope and Target Questions

This paper does not attempt to derive quantum mechanics from pre-physical principles. Instead, it addresses two more specific and well-defined questions: given the existence of complex amplitudes and decohered branch structure, (a) what physical mechanism produces a detector rate of the form $\lambda\_A = \kappa|\psi\_A|^2$? and (b) why does the observable quantity scale as $|\psi\_A|^2$ rather than any alternative function of ψ_A?

These questions are distinct from, and more tractable than, the full derivation of the Born rule. The Born rule makes two claims: that measurement outcomes have probabilities, and that those probabilities equal $|\psi\_A|^2$. The first claim is addressed by the first-passage framework of Taylor (2025a). The second — why $|\psi\_A|^2$ specifically — is the question of this paper.

The answer pursued here does not invoke probability axioms, decision theory, symmetry arguments applied to probability assignments, or Gleason's theorem. Instead, it is derived from three physical facts about how quantum amplitudes couple to fields and how detectors respond to fields. The result is that $|\psi\_A|^2$ is not a postulate about probability but the unique physical measure arising from linear sourcing, quadratic response, and decoherence.

Scope of the result. This paper establishes a physical realisation result, not an independent derivation of the Born rule's ultimate origin. The $|\psi\_A|^2$ structure is already present in the quantum state $|\Psi\rangle$; the κ-field coupling propagates it to detector transition rates through a sequence of physically grounded steps. The κ-field transmits the amplitude-squared weighting from the state to the measurement apparatus — it does not generate that weighting from more primitive principles. Stating this plainly is important: it correctly scopes the contribution as a mechanistic account of how Born-rule structure reaches macroscopic outcomes.

The present result should therefore be understood as identifying the physical realisation of the quadratic measure within quantum mechanics, rather than deriving its ultimate origin. The structure $|\psi|^2$ is already implicit in the Hilbert-space formalism; this paper shows how that structure becomes physically realised in detector dynamics. It explains why $|\psi|^2$ is the quantity that appears in all physically measurable rates, given the structure of quantum amplitudes and detector interactions — not because it is postulated as a probability, but because it is the only form consistent with linear sourcing, quadratic response, and decoherence. The present work therefore identifies the physical origin of the quadratic measure within quantum mechanics, in the sense of how it is realised in detector interactions, while leaving its deeper mathematical origin within Hilbert space as a separate question. Existing derivations of $|\psi|^2$ — Gleason-type results, decision-theoretic approaches, envariance — focus on consistency conditions on probability measures; the present work instead identifies the dynamical mechanism by which this structure becomes physically realised in detector transition rates.

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2. Setup: Post-Decoherence State and κ-Field

After decoherence, the joint state takes the form

$|\Psi\rangle = \Sigma_A \psi_A |A\rangle|E_A\rangle, \langle E_A|E_B\rangle \approx 0.$

The branches are dynamically decoupled. Any measurable quantity must be a functional of the branch amplitudes $\{\psi_A\}$. The question is which functional — a question this paper answers by tracing how $\{\psi_A\}$ propagates through the κ-field to a detector.

In VERSF, the κ-field $\kappa(x,t)$ is sourced by the commitment current $J(x,t)$. Prior to commitment, the detector is embedded in a κ-bath driven by the uncommitted state. The κ-field obeys

$(\partial^2_t - \nabla^2 + m^2_\kappa)\, \kappa(x,t) = \int K(t - t')\, J(x,t')\, dt',$

with memory kernel $K(\tau) \sim \cos(m_\kappa \tau)/\tau$.

Framework postulate: the global κ-field. A foundational feature of VERSF's architecture — not derivable from standard quantum mechanics or decoherence alone — is that the κ-field is defined as a *global* field over configuration space, sourced by all pre-committed branches simultaneously. Each branch contributes to the same κ-field in proportion to its amplitude. This is not the only possible architecture: one could instead assign an independent κ-field sector to each decohered branch, with no global coupling. The derivation of $S_X(\omega_{cm}) \propto \Sigma_A |\psi_A|^2 \tilde{S}_A$ in Section 6 depends on the global structure: cross-branch power terms are present in the initial correlator $\langle X^2\rangle$ and then eliminated by decoherence, which requires all branches to have been coupled to a common field before commitment. This global postulate should be understood as load-bearing in the derivation.

The role of m_κ. The κ-field mass m_κ sets the characteristic frequency of the memory kernel and the bandwidth of κ-field fluctuations driving detector activation. While not derived here,

m_κ is constrained from two directions. From above: m_κ must lie above detector transition frequencies to ensure spectral smoothness at ω_{cm}, while remaining below scales at which κ-field fluctuations would disrupt observed macroscopic stability. From below: the Wigner–Weisskopf pole approximation in Section 5.2 requires $S_X(\omega)$ to be smooth near the detector transition frequency ω_{cm}. Since the memory kernel $K(\tau) \sim \cos(m_\kappa \tau)/\tau$ has spectral features at frequencies set by m_κ, this smoothness condition is satisfied when $m_\kappa \gg \omega_{cm}$ — i.e., when the κ-field's characteristic frequency is well above the detector's transition frequency, placing the memory-kernel spectral structure far from the pole. If $m_\kappa \sim \omega_{cm}$, resonance effects could invalidate the pole approximation and introduce corrections beyond those listed in Section 11. The physically motivated regime for laboratory detectors is therefore $m_\kappa \gg \omega_{cm}$, which places m_κ in a finite, in-principle-measurable window above the GHz–THz range of typical quantum detector transitions. An upper bound on m_κ from VERSF first principles — from cosmological considerations, from the requirement that κ-field fluctuations not disrupt macroscopic classical behaviour at larger scales, or from the κ-field propagator's renormalisation structure — has not yet been derived and is identified as an open problem in Section 14.

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3. Linear Branch-Source Decomposition

The commitment current $J(x,t)$ is an operator on the combined system-environment Hilbert space. Its two-point correlation function in the post-decoherence state is:

$$\langle\Psi|J(x,t)\, J(x',t')|\Psi\rangle = \Sigma_{\{A,B\}}\, \psi^*_A\, \psi_B\, \langle A|\langle E_A|\, J(x,t)\, J(x',t')\, |B\rangle|E_B\rangle.$$

Einselection suppresses the off-diagonal matrix elements of local operators between orthogonal pointer states:

$$\langle E_A|\, O^{env}\, |E_B\rangle \approx \delta_{AB}\, \langle E_A|\, O^{env}\, |E_A\rangle.$$

The two-point function therefore diagonalises:

$$\boldsymbol{\langle\Psi|J(x,t)\, J(x',t')|\Psi\rangle \approx \Sigma_A\, |\psi_A|^2\, \langle A|\langle E_A|\, J(x,t)\, J(x',t')\, |A\rangle|E_A\rangle.}$$

The $|\psi_A|^2$ weighting appears at the level of observable correlation functions. Both the one-point mean $\langle J\rangle$ and the connected two-point function $\langle \delta J\, \delta J\rangle$ diagonalise to $\Sigma_A |\psi_A|^2$ by the same einselection argument.

Operator-level notation. For bookkeeping purposes, write the effective source as:

$$J_{eff}(x,t) = \Sigma_A\, \psi_A\, \xi_A(x,t),$$

where ξ_A is the branch-A component of the current operator. This decomposition reflects the projection of J onto the branch basis defined by $|\Psi\rangle$. The amplitudes ψ_A are state-dependent weights — not intrinsic properties of the operator J — and ψ_A enters linearly as a direct consequence of two structural facts: the quantum superposition principle (the state $|\Psi\rangle$ is a linear combination of branch contributions) and the linearity of the field equations (J is sourced linearly by $|\Psi\rangle$). The amplitude cannot appear squared at this level without violating one of these structural features. The squaring occurs at the expectation-value step, not the sourcing step, exactly as in standard quantum mechanics where $\langle O\rangle = \langle\psi|O|\psi\rangle$ involves $|\psi\rangle$ linearly while yielding bilinear observable quantities.

The linearity is structurally forced: quantum states are superpositions; field equations are linear; therefore sourcing is linear. The amplitude ψ_A enters at first power. The squaring has not yet occurred.

The detector signal is the linear functional of κ:

$$X(t) = \int d^3x\, w(x)\, \kappa(x,t) = \int dt'\, G(t - t')\, J_{eff}(t'),$$

where G is the retarded detector-weighted Green's function. At the operator level, X(t) is linear in ψ_A.

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4. Quadratic Detector Response: Why Detectors Measure Power

No physical detector, in the absence of an external phase reference, measures field amplitude directly. Every detector measures a quantity that is quadratic in the field:

- A photodiode measures photon absorption rate $\propto |E|^2$
- A Geiger–Müller tube measures ionisation density $\propto |E|^2$
- A superconducting nanowire measures deposited energy $\propto X^2$
- A bolometer measures incident power $\propto |B|^2$

This is not a quantum mechanical fact. It is a classical fact about energy transfer: energy flux is always quadratic in field amplitude. For any wave equation — acoustic, electromagnetic, or κ-field — the energy density is quadratic in both field components by the structure of the wave Hamiltonian. For an isolated detector coupled directly to a field bath — the regime relevant here — the leading-order response is quadratic in the field, and this quadratic response is what the FGR transition rate captures. Linear response can be engineered in specific contexts (homodyne detection with a strong local oscillator, certain low-field photoconductive regimes), but these require an external reference amplitude that is not available in the setting considered here, where the κ-field bath is the sole drive. At leading order in the field coupling, the squaring at the detector is universal.

In VERSF's framework, the physically relevant detector quantity is the power spectral density of X(t):

$$S_X(\omega) = \int d\tau\, e^{i\omega \tau}\, \langle X(t)X(t+\tau)\rangle.$$

This is a *second-order* correlation function of the driving field — bilinear in X and therefore, via the linear relation $X \propto J_{eff} \propto \Sigma_A \psi_A \xi_A$, quadratic in the branch amplitudes. The step from first-power sourcing to second-power observable is made here, by the physics of energy detection.

The appearance of $|\psi|^2$ in measurement outcomes is therefore not fundamentally a statement about probability, but about energy transfer. Probability emerges only as a statistical description of first-passage events driven by this energy-coupled interaction — via the competition formalised in Taylor (2025a).

Since $X(t) = \Sigma_A \psi_A Y_A(t)$ (where Y_A is the Green's-function-convolved branch source), the correlator expands as:

$$\langle X(t)X(t+\tau)\rangle = \Sigma_{\{A,B\}} \psi_A \psi_B \langle Y_A(t) Y_B(t+\tau)\rangle.$$

This is now bilinear in the amplitudes. The cross-terms ($A \neq B$) remain to be handled — that is the task of Section 6. But the shift from exponent 1 to exponent 2 has already occurred here, as a direct consequence of the detector measuring power rather than amplitude.

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5. From κ-Field Coupling to Transition Rate: FGR Derivation

The primary derivation uses time-dependent perturbation theory. The earlier Kramers argument employed in a prior draft is withdrawn; see Section 5.3.

5.1 Detector Model and Coupling

The detector is a quantum system with metastable state $|m\rangle$ and committed state $|c\rangle$, separated by transition frequency $\omega_{cm} = (E_c - E_m)/\hbar$. This is the frequency at which the spectral density of the driving field is evaluated — playing the same role as the barrier frequency ω^* in the Kramers–Grote–Hynes framework. Both locate the relevant spectral weight of the drive; they differ in physical interpretation (quantum transition frequency vs classical barrier frequency) but not in their function. The κ-field couples to the detector through:

$$H_{int} = g \cdot \hat{O}_{det} \cdot X(t),$$

where $\hat{O}_{det}$ is the detector's coupling operator and g is the coupling constant.

5.2 Fermi's Golden Rule in the Wigner–Weisskopf Regime

In the weak-coupling regime ($g \ll \Delta E/\hbar$), time-dependent perturbation theory gives the transition rate from $|m\rangle$ to $|c\rangle$:

$\Gamma = (g^2/\hbar^2)\, |\langle c|\hat{O}_{det}|m\rangle|^2 \cdot S_X(\omega_{cm})$,

where $S_X(\omega_{cm})$ is the power spectral density at the detector transition frequency (Cohen-Tannoudji et al., 1992, ch. IV). The rate is proportional to S_X — a second-order spectral quantity, bilinear in the source, for all noise spectra including non-Markovian.

Validity: weak coupling. FGR requires $\Gamma \ll \omega_{cm}$. For single-quantum-sensitive detectors, this is satisfied in the pre-commitment regime of tick production.

Irreversibility and the Wigner–Weisskopf regime. The committed state $|c\rangle$ is not a stationary eigenstate — the detector escapes into it irreversibly. Standard FGR assumes non-depletion of the initial state and reversible transitions, which would seem to disqualify it here. The correct generalisation is the Wigner–Weisskopf approximation: irreversible decay from $|m\rangle$ into a continuum of committed states $\{|c_f\rangle\}$, analogous to spontaneous emission. For this approximation to hold, the detector's committed-state manifold must be sufficiently dense near ω_{cm} to function as an effective continuum. For macroscopic single-quantum detectors — photodiodes, SNSPDs, Geiger tubes — the density of accessible final states (electronic, phononic, thermal modes) near the relevant energy scale typically far exceeds the coupling bandwidth, making the continuum approximation excellent in practice. In this regime the decay rate is still $\Gamma = (g^2/\hbar^2)|\langle c|\hat{O}_{det}|m\rangle|^2\, S_X(\omega_{cm})$ summed over final states — the FGR formula unchanged in form — but the initial-state population decays exponentially as $e^{-\Gamma t}$ and the transition is genuinely irreversible. The VERSF detector operates precisely in the Wigner–Weisskopf regime.

Non-Markovian κ-field and the Wigner–Weisskopf regime. The κ-field memory kernel $K(\tau) \sim \cos(m_\kappa \tau)/\tau$ produces algebraically decaying correlations with spectral structure concentrated near frequencies of order m_κ. A potential tension exists with the Wigner–Weisskopf result, which in its standard form assumes exponential (Markovian) decay. The resolution requires specifying a condition rather than merely asserting one. The Wigner–Weisskopf pole approximation holds when $S_X(\omega)$ is sufficiently smooth in the neighbourhood of ω_{cm} — but whether this is satisfied depends on the relation between m_κ and ω_{cm}. If $m_\kappa \gg \omega_{cm}$, the spectral features of the memory kernel lie far above the detector transition frequency, $S_X(\omega)$ is smooth near ω_{cm}, and the pole approximation is valid. If $m_\kappa \sim \omega_{cm}$, resonance effects could introduce structure in S_X near the pole and invalidate the approximation, producing corrections beyond those in Section 11. The physically motivated regime for laboratory detectors — where ω_{cm} is in the GHz–THz range — therefore requires $m_\kappa \gg \omega_{cm}$, which is the lower bound on m_κ identified in Section 2. In this regime, the non-Markovian memory produces corrections to early-time dynamics — transient oscillations and power-law tails — but does not alter the leading-order exponential rate or the $|\psi_A|^2$ functional dependence, which depends only on $S_X(\omega_{cm})$, not on the shape of S_X away from ω_{cm}. Outside this regime, deviations from exponential decay are expected, but they modify the prefactor κ rather than the $|\psi_A|^2$ form. Both the FGR (microscopic) and Grote–Hynes (coarse-grained) descriptions are therefore controlled by the same spectral density $S_X(\omega)$; non-Markovian memory modifies early-time dynamics and the prefactor κ but not the $|\psi_A|^2$ functional dependence, which depends only on $S_X(\omega_{cm})$. The two frameworks are not in tension — they are two descriptions of the same spectral physics at different scales.

5.3 Withdrawal of the Kramers Argument and Status of the Grote–Hynes Result

The standard overdamped Kramers rate $\lambda \sim (\omega_0\omega_b/2\pi\gamma) \cdot \exp(-\Delta U/D)$ is *exponentially* dependent on noise power D — not linearly proportional to it. An earlier formulation of this derivation claimed $\lambda \propto \langle X^2\rangle$ from Kramers theory; this was incorrect in the generic regime and is withdrawn.

The Grote–Hynes generalisation of Kramers theory for colored noise confirms that the rate depends on $S_X(\omega^*)$ at the barrier frequency — a bilinear spectral quantity — which independently supports the $|\psi_A|^2$ form. The FGR derivation applies at the microscopic quantum level; the Kramers–Grote–Hynes framework describes coarse-grained classical escape over a barrier. The agreement in functional dependence reflects that both are controlled by the same spectral density $S_X(\omega)$ evaluated at the relevant dynamical frequency (ω_{cm} for FGR, ω^* for the barrier). **Grote–Hynes is cited here as corroborating evidence, not as a load-bearing component of the derivation.** The FGR argument in Section 5.2 is self-contained; Grote–Hynes provides independent confirmation at a different description level that the bilinear spectral dependence is robust under non-Markovian memory. The primary derivation is FGR.

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6. Decoherence ⇒ Diagonal Noise Correlator

The power spectral density $S_X(\omega_{cm})$ involves the two-point correlator of X:

$$\langle X(t)X(t+\tau)\rangle = \Sigma_{\{A,B\}}\ \psi_A\ \psi_B\ \langle Y_A(t)\ Y_B(t+\tau)\rangle.$$

The cross-branch correlator $\langle Y_A(t)\ Y^*_B(t+\tau)\rangle$ is proportional to $\langle E_A|\delta J^{env}(t)\ \delta J^{env}(t+\tau)|E_B\rangle$. Since δJ^{env} is a local environmental operator, einselection applies: this matrix element is suppressed by $|\langle E_A|E_B\rangle|$ for $A \neq B$.

Connecting orthogonality to statistical independence. The physical content is: branches that have left distinguishable macroscopic records in the environment cannot coherently source correlated κ-field fluctuations at the detector, because any such correlation would require those records to overlap. Strictly, orthogonality ensures suppression of cross-correlators to the same order as interference terms; statistical independence is therefore an approximation valid when environmental records are sufficiently redundant and local observables cannot access global coherence. Therefore:

$$\langle Y_A(t)\ Y^*_B(t+\tau)\rangle \approx \delta_{AB}\ C_A(\tau),$$

where $C_A(\tau)$ is the within-branch autocorrelation. The noise correlator diagonalises:

$$\langle X(t)X(t+\tau)\rangle \approx \Sigma_A\ |\psi_A|^2\ C_A(\tau),$$

and the spectral density is:

$S_X(\omega_{cm}) \approx \Sigma_A |\psi_A|^2 \tilde{S}_A(\omega_{cm})$,

where $\tilde{S}_A(\omega_{cm}) = \int d\tau\, e^{(i\omega_{cm} \tau)} \iint G\, G\, C_A$ is the per-branch spectral weight.

The three physical ingredients have now acted in sequence. Linear sourcing gave $X \propto \Sigma_A \psi_A Y_A$ (amplitude enters at first power). Quadratic detector response gave $S_X \propto \langle X^2 \rangle$ (exponent shifts to 2). Decoherence diagonalisation gave $S_X \approx \Sigma_A |\psi_A|^2 \tilde{S}_A$ (cross-terms vanish). The $|\psi_A|^2$ has emerged as the unique diagonal residue.

Note on the foundational status of decoherence. This step assumes $\langle E_A|E_B \rangle \approx 0$ as established by unitary evolution. Characterising decoherence quantitatively typically uses the quantum formalism in ways that may implicitly involve amplitude-squared weighting for state norms. The derivation therefore constitutes a partial reduction: given decoherence, the $|\psi_A|^2$ weighting follows. A fully foundational derivation — establishing decoherence without Born-rule assumptions — remains an open problem shared by all current approaches in this programme (Zurek 2003; Wallace 2012; Deutsch 1999).

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7. Branch-Blindness ⇒ Common Prefactor and Error Estimate

From the FGR expression, the per-branch prefactor is:

$\kappa_A = (g^2/\hbar^2) |\langle c|\hat{O}_{det}|m \rangle|^2 \tilde{S}_A(\omega_{cm})$.

The matrix element $|\langle c|\hat{O}_{det}|m \rangle|^2$ is determined entirely by the detector's internal structure and is branch-independent. Variations in κ_A arise solely from differences in $\tilde{S}_A(\omega_{cm})$ — the κ-field spectral density within each branch's sector.

After decoherence, local κ-field fluctuations within each branch arise from the same underlying thermal and vacuum bath modes. Branch-specific corrections to $C_A(\tau)$ come from the branch's own back-action on the κ-bath, an effect of order $|\psi_A|^2$ relative to the background:

$\kappa_A = \kappa (1 + \delta_A)$, $|\delta_A| \sim |\psi_A|^2 \cdot N_{bath}^{-1}$,

where $N_{bath} \sim 10^{23}$ for a macroscopic detector. The resulting deviation from Born statistics:

$|P(A) - |\psi_A|^2| \sim |\psi_A|^4 \cdot N_{bath}^{-1} \sim 10^{-23}$,

is unmeasurably small, consistent with the high-precision Born rule agreement observed experimentally.

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8. The Rate Law: $\lambda_A = \kappa|\psi_A|^2$

Substituting into the FGR rate formula:

$\Gamma \propto S_X(\omega_{cm}) \approx \Sigma_A |\psi_A|^2 \tilde{S}_A(\omega_{cm})$.

With $\kappa_A \approx \kappa$, the branch-resolved first-passage rate is:

$\lambda_A = \kappa |\psi_A|^2$.

This is the rate law. The derivation chain to this point is:

1. State $|\Psi\rangle$ is linear in $\psi_A \rightarrow J_{eff} = \Sigma_A \psi_A \xi_A$ (linear sourcing, forced)
2. Detector measures power, not amplitude → rate $\propto S_X(\omega_{cm})$ (quadratic response, universal)
3. Decoherence → $S_X \approx \Sigma_A |\psi_A|^2 \tilde{S}_A$ (cross-terms vanish)
4. Branch-blindness → $\tilde{S}_A \approx \kappa$ (corrections $\sim N_{bath}^{-1}$)
5. Result: $\lambda_A = \kappa|\psi_A|^2$

The $|\psi_A|^2$ did not come from a probability postulate. It came from the interaction between linear quantum sourcing and the universal physics of quadratic detector response, with decoherence providing the diagonal structure.

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9. Uniqueness: Why Not $|\psi_A|^4$ or Any Other Form?

The derivation above shows that $|\psi_A|^2$ arises at leading perturbative order. The uniqueness argument maps out precisely how the constraint set narrows the space of possible rate functionals — and where each constraint does and does not do work. Four constraints apply to the per-branch rate functional $f(\psi_A)$, and it is important to be precise about what each one achieves:

Constraint 1 — Locality. Rates depend on local field interactions only; no action-at-a-distance in the detector-field coupling.

Constraint 2 — Analyticity. f is a smooth function of ψ_A. A non-analytic rate response to a continuously varying amplitude is unphysical for any real detector.

Constraint 3 — U(1) phase covariance. The global phase $\psi_A \rightarrow e^{i\theta} \psi_A$ is physically unobservable; f must satisfy $f(e^{i\theta} \psi_A) = f(\psi_A)$. This eliminates all terms of the form $\psi_A^n (\psi_A^*)^m$ with $n \neq m$, leaving only terms of the form $c_n |\psi_A|^{2n}$.

Constraint 4 — Additivity. For independent branches A and B — whose noise processes are statistically uncorrelated by decoherence — the combined rate is $f(\psi_A) + f(\psi_B)$. Independent sources add without cross-terms.

What the four constraints together achieve: Constraints 1–3 restrict f to the polynomial ring of terms $c_n|\psi_A|^{2n}$ — an infinite family. Constraint 4 (additivity) is satisfied by every term in this family independently; it does not by itself exclude $|\psi_A|^4$ or higher powers. The actual selection of $|\psi_A|^2$ as the leading term requires a fifth input: the perturbative order of the FGR calculation. This is where the work is done.

Examining the fate of each candidate under this two-stage analysis:

- **$|\psi_A|^2$ (n = 1):** admitted by all four constraints. This is the FGR leading-order term, proportional to the second-order correlator $S_X(\omega_{cm})$, arising at order g^2 in the coupling.
- **$|\psi_A|^4$ (n = 2) and higher:** admitted by all four constraints — including additivity, which these terms satisfy independently across branches. They are excluded by perturbative order: $|\psi_A|^4$ corresponds to the fourth-order correlator of the source, arising at order g^4, suppressed by $(g/\Delta E)^2$ relative to the leading FGR term. More precisely: $f(\psi_A) = \kappa|\psi_A|^2 + \kappa_2|\psi_A|^4 + \ldots$ is a valid polynomial satisfying all four constraints, but $\kappa_2/\kappa \sim (g/\Delta E)^2$, the weak-coupling parameter that defines the regime of validity. These corrections are present but parametrically small, calculable, and in principle detectable via the $k > 1$ prediction.
- **Non-analytic exponents $|\psi_A|^\alpha$, $\alpha \notin 2\mathbb{Z}$:** these cannot be generated by any finite-order local field interaction. They require infinite-order perturbation theory or non-local couplings, excluded by constraints 1 and 2.
- **$|\psi_A|$ ($\alpha = 1$):** excluded by phase covariance, since $|\psi_A|$ cannot be expressed as a finite polynomial in ψ_A and ψ_A^*, in violation of the analyticity constraint (Constraint 2). It is not a member of the polynomial ring $\mathbb{C}[\psi_A, \psi_A^*]$ at any finite order.

The unique leading-order local analytic rate functional within the perturbative expansion consistent with U(1) covariance is $|\psi_A|^2$. Non-quadratic forms do not arise at any fixed order in a local analytic expansion consistent with U(1) covariance; deviations require higher-order terms that lie beyond the weak-coupling regime in which the derivation is valid. Higher-order terms present at higher perturbative orders are suppressed by $(g/\Delta E)^{2(n-1)}$ and are in principle testable via multi-trigger detectors.

The quadratic form $|\psi_A|^2$ is not a postulate of probability theory. It is the unique physical measure arising from the interaction between linear quantum amplitudes and quadratic detector response under decoherence, at leading perturbative order. Unlike Gleason-type results, which establish $|\psi|^2$ as a consistency condition on probability measures over Hilbert-space subspaces, the present derivation identifies $|\psi|^2$ as a consequence of physical interaction between quantum amplitudes and detector dynamics — a dynamical result, not a measure-theoretic one.

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10. First-Passage Selection and the k > 1 Prediction

Given $\lambda_A = \kappa|\psi_A|^2$ and proportional hazards $h_A(t) = \lambda_A h_0(t)$, the first-passage selection law (derived in Taylor 2025a) yields:

$P(A) = \lambda\_A / \Sigma\_B\ \lambda\_B = |\psi\_A|^2 / \Sigma\_B\ |\psi\_B|^2$.

For normalised states, **$P(A) = |\psi\_A|^2$**. The Born rule follows from κ-field driven threshold dynamics.

The k > 1 prediction. The most direct falsifiable consequence of the rate law is the deviation from Born statistics for detectors requiring multiple independent threshold events. For k = 2, the first-passage probability for branch 1 against branch 2 via competing Gamma(2, λ_A) processes is:

$P_1(k=2) = \lambda_1^2(\lambda_1 + 3\lambda_2) / (\lambda_1 + \lambda_2)^3$.

For $\lambda_1 = 2\lambda_2$ — corresponding to $|\psi_1|^2 = 2|\psi_2|^2$ — this evaluates to 20/27 ≈ 0.741, compared to the k = 1 Born rule prediction of 2/3 ≈ 0.667. The full derivation via the Gamma competition integral is given in Taylor (2025a), Section 8.2.

This prediction is not a modification of quantum theory at the level of states or Hamiltonians, but a modification at the level of detector dynamics. It therefore provides a direct experimental test of whether the Born rule arises from fundamental postulates or from physical threshold processes — a distinction that no existing interpretation of quantum mechanics makes testable in this way. The prediction is quantitative and model-complete: for a specified k and {λ_A}, the deviation from Born statistics is fixed by the Gamma competition formula, independent of interpretational assumptions. This prediction is therefore not interpretational but dynamical: it distinguishes between theories at the level of detector physics, not state description.

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11. Domain of Validity and Corrections

The derivation relies on:

- Decoherence: ⟨E_A|E_B⟩ ≈ 0, taken as established — partial reduction (Section 6 and 12).
- Weak coupling: FGR requires $\Gamma \ll \omega\_cm$; satisfied in the pre-commitment linear-response regime.
- Wigner–Weisskopf smoothness: $S\_X(\omega)$ smooth near ω_cm so the pole approximation holds.
- Branch-blind detector: κ_A variations suppressed by N_bath^{-1}.
- Metastable, single-trigger regime (k = 1): thermodynamically forced for single-quantum-sensitive detectors.
- Analytic rate functional: uniqueness of $|\psi\_A|^2$ holds at leading order in the polynomial expansion.

Calculable corrections:

Source	Effect	Order of magnitude
Residual cross-branch correlations	Non-diagonal terms ∝ $\psi_A \psi^*_B$	~ $|\langle E_A|E_B\rangle|$
Branch-dependent κ_A	Deviations from exact $|\psi|^2$	~ N_{bath}^{-1}
Higher-order FGR (g^4 terms)	Corrections to λ_A	~ $(g/\Delta E)^2$ × leading order
Multi-trigger (k>1)	Gamma-distributed waiting times	Taylor (2025a), Section 8.2
Non-Gaussian noise	Higher-cumulant corrections	~ (fluctuation amplitude / $\Delta E)^2$
Non-Markovian memory	Shifts spectral weight at ω_{cm}	Absorbed into κ, no form change

12. What This Derivation Establishes and Does Not Establish

What this derivation does NOT do:

- It does not derive Hilbert space, complex amplitudes, or quantum superposition from pre-quantum principles. The present result should therefore be understood as identifying the physical realisation of the quadratic measure *within* quantum mechanics, rather than deriving its ultimate origin. This deeper question — why $|\psi|^2$ is built into the Hilbert-space formalism in the first place — is addressed within the VERSF programme by the physical necessity results of Taylor (2025, *Physical Necessity of Quantum Probability Structure*), which derives the admissibility of $|\psi|^2$ from pre-quantum principles of finite distinguishability, irreversible commitment, and temporal extensibility.
- It does not derive decoherence. The condition $\langle E_A|E_B\rangle \approx 0$ is taken as established by unitary evolution.
- It constitutes a partial reduction conditional on decoherence: given $\langle E_A|E_B\rangle \approx 0$, $|\psi_A|^2$ is the unique leading-order physically realised measure. A ground-up derivation establishing decoherence without amplitude-squared assumptions — which would close the remaining circularity — is a separate open problem shared by all current approaches in the decoherence programme.
- It does not prove the rate law non-perturbatively. The uniqueness of $|\psi_A|^2$ holds at leading order in weak coupling; higher-order corrections are suppressed but present.

What this derivation DOES establish:

- That once complex quantum amplitudes and decoherence are granted, $|\psi_A|^2$ is the unique physically realised measure at leading perturbative order.
- That the squaring arises from a universal physical fact — detectors measure power, not amplitude — and not from a probability axiom.
- That decoherence is the mechanism by which cross-branch interference is eliminated, leaving the diagonal $|\psi_A|^2$ structure.

- That no alternative functional form survives all four physical constraints at leading order.
- That the result is not specific to VERSF's κ-field: any quantum framework in which branch amplitudes source an intermediary field linearly, detectors couple quadratically to that field, and branches decohere must produce $|\psi_A|^2$ weighting by the same three-step argument. The κ-field is VERSF's specific mediator. Whether detector frameworks without an explicit intermediary field — where quantum branches couple directly to the detector without a mediating classical background — produce the same result is a separate question not addressed here; the present derivation requires the intermediate-field structure as a load-bearing element.

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13. The Complete Derivation Chain

This paper is the second in a two-paper programme. Together with Taylor (2025a), the complete chain from quantum amplitudes to measurement probabilities runs as follows:

This paper establishes: quantum amplitude → κ-field sourcing (linear, forced) → detector response (quadratic, universal) → decoherence (diagonal) → rate law $\lambda_A = \kappa|\psi_A|^2$ (uniquely, at leading order).

Taylor (2025a) establishes: given $\lambda_A = \kappa|\psi_A|^2$, proportional hazards, and first-passage competition → $P(A) = |\psi_A|^2$ (Born rule as theorem of competing Poisson processes).

The full chain:

Quantum amplitude → κ-field sourcing (linear) → power spectral density (quadratic) → decoherence (diagonal) → rate law $\lambda_A = \kappa|\psi_A|^2$ → first-passage competition → Born rule $P(A) = |\psi_A|^2$.

At each step, the physical mechanism is identified, the approximations are stated, and the validity conditions are specified. No step invokes the Born rule itself. The chain is grounded in field physics, perturbation theory, and survival statistics.

The essential shift introduced across this programme: quantum probabilities are not properties of static quantum states alone, but of the dynamical coupling between those states and irreversible threshold processes. The Born rule is a statement about physical interaction, not merely about representation.

Within the VERSF programme, the present result should be read as the intermediate layer in a three-part structure: the foundational derivation of the admissibility of $|\psi|^2$ (Taylor 2025, *Physical Necessity of Quantum Probability Structure*), the present derivation of its physical realisation in detector dynamics, and the companion derivation of outcome probabilities via first-passage competition (Taylor 2025a). Each layer is independently grounded; together they constitute a complete and non-circular account of quantum measurement probability.

14. Open Problems

Non-perturbative computation of κ. The prefactor $\kappa = (g^2/\hbar^2)|\langle c|\hat{O}_{det}|m\rangle|^2 \tilde{S}_{bath}(\omega_{cm})$ depends on the full κ-field spectral density, involving the memory kernel $K(\tau)$. A non-perturbative evaluation would yield a concrete prediction for κ as a function of detector parameters and m_κ.

Foundational status of decoherence. The derivation assumes decoherence as established. Closing this loop — deriving $\langle E_A|E_B\rangle \approx 0$ without amplitude-squared weighting — is the residual circularity shared by all current approaches in the decoherence programme.

Extension to entangled systems. For entangled systems measured by spatially separated detectors, the joint branch structure requires a configuration-space generalisation of the noise correlator diagonalisation. The einselection argument is expected to go through for joint environmental orthogonality, but deserves explicit treatment.

Rigorous correction bounds. The N_{bath}^{-1} and $(g/\Delta E)^2$ estimates are order-of-magnitude arguments. Rigorous bounds require specifying κ-field mode density, coupling bandwidth, and detector response profile — tractable within VERSF's field theory.

Self-consistency of the classical stochastic treatment of X(t). The detector coupling $H_{int} = g \cdot \hat{O}_{det} \cdot X(t)$ treats $X(t)$ as a classical stochastic background — a c-number drive sourced by the quantum branches. Establishing that this treatment is self-consistent with the full quantum κ-field requires verifying that quantum fluctuations of κ are sub-leading relative to classical power in the pre-commitment regime. This requires a one-loop calculation of the κ-field propagator within VERSF's field theory and is deferred to future work.

15. Summary

Why does anything physical scale as $|\psi_A|^2$? Three physical facts act together:

Fact 1 — Linear sourcing (forced). Quantum states are linear objects; field equations are linear operators; therefore the κ-field is sourced with ψ_A at first power. The amplitude enters linearly. The squaring has not yet occurred.

Fact 2 — Quadratic detector response (universal). No physical detector measures field amplitude directly. Every detector measures power — energy, intensity, flux — which is always a second-order quantity in the field. The FGR transition rate is proportional to the spectral density $S_X(\omega_{cm})$, a second-order correlator. This is where the exponent shifts from 1 to 2.

Fact 3 — Decoherence diagonalisation (enforced by branch distinctness). After decoherence, branches have left distinguishable environmental records, suppressing cross-branch noise correlations. The spectral density diagonalises to $\Sigma_A |\psi_A|^2 \tilde{S}_A$. The cross-terms vanish; the $|\psi_A|^2$ is the unique diagonal residue.

Why the exponent is 2 and nothing else: uniqueness follows from four constraints — locality, analyticity, U(1) covariance, and additivity — which admit only $|\psi_A|^2$ from the analytic expansion of any local rate functional at leading perturbative order. Higher even powers are suppressed by the weak-coupling parameter; non-analytic forms violate locality.

The central conclusion: The quadratic form $|\psi|^2$ is not a postulate of probability theory but the unique physical measure arising from the interaction between linear quantum amplitudes and quadratic detector response under decoherence. It is a property of physical interaction, not of representation. The experimental consequence — that engineered detectors with $k > 1$ threshold requirements will deviate from Born statistics by a fixed, calculable amount — is the direct empirical shadow of this physical origin. The present result therefore identifies the physical origin of the quadratic measure in measurement processes, while complementary VERSF results address its deeper mathematical necessity.

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16. Anticipated Objections and Responses

Objection 1: Circularity — you assume $|\psi|^2$ to get $|\psi|^2$.

Response. The derivation does not assume $|\psi_A|^2$ is a probability or a rate. It derives a rate functional from field physics: linear sourcing, quadratic power, decoherence diagonalisation. The $|\psi_A|^2$ emerges as the residue of these three physical facts. The paper establishes a structural reduction: the measurement problem is reduced to a dynamical question about κ-field coupling, rather than solved from scratch. Completion at a fully non-perturbative level, and the elimination of the decoherence precondition, are identified as open problems.

Objection 2: The "quadratic response" claim restates that detectors measure energy — but classical energy ∝ $|E|^2$ is itself tied to the Born rule via quantum field theory.

Response. The quadratic dependence of energy on field amplitude is a classical field-theoretic fact, independent of quantum mechanics. For any wave equation — acoustic, electromagnetic, or κ-field — energy density is quadratic in the field by dimensional analysis and the structure of the wave Hamiltonian. This is true independently of the Born rule. Furthermore, at the level of the detector-κ-field coupling, X(t) is treated as a *classical stochastic background* — a c-number drive, not a quantised field operator. In this regime the Born rule for quantised fields is simply not in play, because there is no second quantisation at this coupling step. The connection between quantum amplitudes and κ-field energy is made by $H_{int} = g \cdot \hat{O}_{det} \cdot X(t)$, where X(t) is the classical stochastic signal sourced by the quantum branches. This classical stochastic treatment of X(t) is a modelling assumption of the framework at leading order — it is the standard approximation used when a quantum system is driven by a field whose quantum fluctuations are sub-leading relative to its classical power. Establishing the full self-consistency of this treatment with the quantum κ-field structure is an open problem, identified as such in Section 14.

Objection 3: Proportional hazards is an ad hoc assumption.

Response. The factorisation $h_A(t) = \lambda_A h_0(t)$ reflects a physical separation between two distinct processes. $h_0(t)$ encodes the detector's internal relaxation and amplification dynamics, determined entirely by the detector's construction and independent of the quantum branch driving it. The branch label enters only through the strength of the κ-field drive, setting the scale λ_A. Deviations from exact proportional hazards produce calculable corrections to Born statistics (Section 11).

Objection 4: How does a single detector couple to multiple branches after decoherence?

Response. The detector does not couple to multiple independent systems. It couples to a single κ-field defined over configuration space, encoding contributions from all pre-committed branches weighted by their amplitudes. Decoherence suppresses cross-branch interference at the level of local observables but does not eliminate each branch's contribution to the global κ-field. Outcome selection is competition between contributions to a single stochastic driving field.

Objection 5: The FGR derivation may not apply to an irreversible threshold process.

Response. Addressed in Section 5.2. The correct framework is the Wigner–Weisskopf approximation — irreversible decay of the metastable state $|m\rangle$ into a continuum of committed states. In this regime the rate formula $\Gamma = (g^2/\hbar^2)|\langle c|\hat{O}_{det}|m\rangle|^2 S_X(\omega_{cm})$ is unchanged in form; the difference is that the initial-state population decays exponentially and the transition is irreversible. The $|\psi_A|^2$ dependence follows by the same argument.

Objection 6: The uniqueness argument at leading perturbative order is weak — the true exponent could differ from 2 at higher orders.

Response. Correct, and the paper does not overclaim. The uniqueness of $|\psi_A|^2$ holds at leading order in weak coupling. Higher-order corrections are suppressed by $(g/\Delta E)^{2(n-1)}$ and are in principle detectable via multi-trigger detectors — the $k > 1$ prediction. The claim is not that $|\psi_A|^2$ holds exactly for all couplings, but that it is uniquely selected at leading order, with deviations calculable and testable.

Objection 7: The decoherence step is circular.

Response. Acknowledged in Section 12 and Section 6. Whether establishing decoherence quantitatively requires amplitude-squared weighting is a foundational question the present derivation does not resolve. The paper constitutes a partial reduction: given decoherence as established, $|\psi_A|^2$ is the unique leading-order measure. This limitation is stated explicitly rather than concealed.

Objection 8: The predictions may not be experimentally accessible.

Response. The framework makes two quantitative predictions beyond standard quantum mechanics. First, multi-trigger detectors ($k > 1$) exhibit calculable deviations from Born statistics — $20/27 \approx 0.741$ versus $2/3 \approx 0.667$ for $k = 2$ and $\lambda_1 = 2\lambda_2$; the statistical requirements are derived in Taylor (2025a). Second, the timing distribution of measurement outcomes is predicted to depend on the total κ-field coupling κ, which is in principle measurable. Both predictions are quantitative and independent of interpretational assumptions.

This paper (Taylor 2025b) derives, at leading order under physically controlled approximations, the rate law and its unique quadratic form. The Born rule derivation enabled by this rate law is carried out in the companion paper Taylor (2025a). Together, these results reduce the quantum measurement problem to a fully dynamical statement: outcome probabilities are determined by the coupling of pre-committed branch structure to κ-field-driven irreversible threshold processes.

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