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RESEARCH

Contextual Bohmian Quantum Field Theories: A Hylomorphic Approach to QFT

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Abstract

We propose an extension of Bell-type Bohmian quantum field theories, called *Contextual Bohmian Quantum Field Theory* (CBQFT), which integrates micro-level dynamics and macro-level contextual structure within a unified, ontologically explicit formalism. CBQFT introduces classical variables Λ that encode macroscopic contexts—such as detector configurations, thermal phases, or symmetry-breaking sectors—and allows these to modulate the underlying quantum dynamics in a lawlike way. We develop two versions of the model. CBQFT-1 treats context as a fixed but dynamically influential background, entering via a context-sensitive Hamiltonian and modified Bell-type jump rates on a single Fock space. CBQFT-2 upgrades context to a *dynamical* variable co-evolving with the particle (or field) configuration: $\Lambda(x,t)$ selects a (typically inequivalent) representation of the field algebra on a Hilbert space, wavefunctions are realised as global sections of the resulting Hilbert bundle, and Bohmian trajectories are guided by globally well-defined velocity fields constructed from local currents. Context transitions in CBQFT-2 are governed by a stochastic kernel informed by particle (or field) configurations and histories. This yields a Bohmian QFT with an explicit feedback loop between quantum events and macroscopic structure, offering a hylomorphic account of measurement, decoherence, and top–down causation.

Keywords Bohmian mechanics · Bohmian quantum field theory · Bell-type quantum field theory · Hylomorphism · Hylomorphic primitive ontology approach

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1 Introduction

Bohmian mechanics (BM), in its original non-relativistic form, is an interpretation of quantum theory that posits a fixed number of particles with definite positions, guided by a wavefunction evolving according to the Schrödinger equation [5, 11]. It is often praised by the philosophically minded for offering a clear and realist physical ontology, in contrast with operationalist or instrumentalist interpretations of quantum mechanics. However, extending BM to relativistic and field-theoretic domains remains non-trivial. In standard quantum field theory (QFT), particle number is not conserved: particles may be created or annihilated as a consequence of field interactions [42].

Bell-type Bohmian quantum field theories (BQFTs) have shown that it is possible to ascribe definite particle trajectories even in the presence of creation and annihilation events [12, 58]. In these models, particles follow well-defined world lines that begin or end at discrete events, thereby providing a beable-based account of quantum field processes. The particle configuration evolves deterministically between jumps, and then stochastically undergoes transitions when creation or annihilation occurs. These extensions retain empirical agreement with standard QFT while furnishing a concrete picture of "what is going on". Nonetheless, the structure of the jump process—its timing, causes, and modulation—remains open to interpretation.

Existing BQFTs treat creation and annihilation as purely bottom-up phenomena: the stochastic jump rates are entirely determined by the microscopic wavefunction dynamics. But in standard quantum physics, it is well known that the macroscopic *context*—such as the presence of a measuring device or the surrounding environment—can influence the quantum behaviour of a microscopic system. For example, the Purcell effect shows that an atom's spontaneous emission rate can be modified by placing it in a resonant cavity, i.e., by altering the electromagnetic environment [43]. Related boundary-sensitive phenomena, such as Casimir forces, in which vacuum stresses depend on macroscopic geometry and on the way vacuum energy is regularised, provide further examples of context-dependent behaviour [24, 52, 56]. This suggests that top-down influences from higher-level structures may modulate micro-level processes. Indeed, some physicists have argued that such context-dependence may play an essential role in quantum theory and should be understood as a form of top-down causation [16].

In this paper, we propose two extensions of Bell-type BQFT that incorporate macroscopic context into the dynamics of particle creation and annihilation; two *Contextual* BQFTs. The first model, CBQFT-1, introduces classical context variables Λ —such as thermodynamic conditions, boundary structures, or detector settings—that modulate the stochastic jump rates through additional Hamiltonian terms. These contextual influences are treated as external parameters that exert top-down causal influence on particle dynamics. The second model, CBQFT-2, goes further by allowing the particle (or field) configuration to influence and modify the context itself. Context is now encoded not merely as a parameter but via a choice of unitarily inequivalent representations of the field algebra, enabling changes in global structure (e.g., phase transitions or symmetry-breaking events) that reshape the Hilbert space. There is feedback: the evolution of the microscopic system can trigger shifts in the macroscopic organisation, completing the causal loop. In idealised conditions—where the context is static or trivial—both models reduce to the standard Bell-type BQFT as a limiting case. We thus preserve empirical adequacy while expanding the interpretive resources available.

Pragmatically, CBQFT-1 functions as a context-parametrised extension of Bell-type BQFT that models environmental and apparatus influence directly at the level of the jump law, rather than by enlarging the quantum system until the context is treated fully quantum-mechanically. Conceptually, CBQFT-2 treats macroscopic context as a dynamical variable that encodes local macroscopic forms and is updated by micro-events via sector transitions. This opens conceptual space for coupling Bohmian QFT to genuinely classical structures—for example, a classical gravitational or cosmological background as part of the context which modifies vacuum structure [32]. The structure of the paper is as follows.

In Section 2 we review Bell-type Bohmian mechanics in the context of QFT, highlighting how particle configurations evolve through deterministic flows punctuated by stochastic jumps. Section 3 introduces our first extension, CBQFT-1: we define "context" in physical and dynamical terms and propose a modified stochastic law. In Section 4 we formalise this model, analyse norm conservation and equivariance, and show how the standard non-contextual theory is recovered when contextual contributions vanish. In Section 5 we explore the implications of allowing macroscopic variables to influence quantum processes, including for measurement, the arrow of time, and effective wave function collapse. Section 6 introduces our second extension, CBQFT-2, which treats context as a dynamical variable, admits unitarily inequivalent representations of the field algebra, and enables feedback from micro-level events to macroscopic structure. In Section 7 we axiomatise this richer theory, derive key properties such as norm conservation and equivariance, and establish consistency with Bohmian guidance. Section 8 explores the metaphysical, causal, and empirical implications of incorporating these context-sensitive representations within Bohmian theory, discussing information flow, measurement, and relativistic constraints, whilst advancing a hylomorphic reading in which physical objects are composites of *both* matter *and* form. Section 9 compares our proposals with the Everett (Many Worlds) interpretation and collapse models (GRW/CSL/CWC). Section 10 concludes.

2 BM in QFT (Bell-Type QFT)

BM was originally developed for nonrelativistic quantum mechanics with a fixed number of particles N. It introduces an actual configuration, $Q(t) = (\mathbf{x}_1(t), \mathbf{x}_2(t), \ldots, \mathbf{x}_N(t))$, and postulates that $Q(t)$ evolves according to a deterministic velocity field derived from a wavefunction $\Psi(q,t)$ that is the solution of Schrödinger's equation [13]. The velocity law ensures that if $Q(t_0)$ is distributed according to $|\Psi(q,t_0)|^2$, then $Q(t)$ remains $|\Psi|^2$-distributed for all time. This requirement concerning the initial particle configuration, known as the *equivariance condition*, ensures the preservation of the Born rule by the dynamics. The particles evolve such that their distribution continues to match the quantum probability density at all times, preserving agreement with standard quantum predictions.

To extend this picture to QFT, where the number of particles is not fixed, the Bell-type QFT approach generalises the configuration space to include sectors of different particle counts [10–12]. Specifically, the state vector lives in Fock space $F = \bigoplus_{n=0}^{\infty} \mathcal{H}_n$, where $\mathcal{H}_n$ is the *n*-particle Hilbert subspace. Here, *n* plays the role of a *sector label* and is allowed to vary in time as particles are created or annihilated, in contrast to the fixed *N* of the Bohmian theory above. A point in the total configuration space *Q* is thus an unordered set of particle positions for some *n*, or equivalently $(n, \mathrm{x}_1, \ldots, \mathrm{x}_n)$, with *n* variable, and the squared amplitude of the wavefunction $|\Psi(q)|^2$ defines a probability density on this entire configuration space, spanning all sectors. In other words, this Bohmian version of QFT associates the quantum state Ψ with a distribution of possible particle configurations which have varying particle number. Let's consider how such a system evolves.

2.1 Deterministic Motion Within Sectors

In the absence of creation or annihilation events, the particle positions evolve along continuous trajectories guided by the wavefunction. In practice, this means that on each *n*-particle sector we have a guiding equation analogous to the usual Bohmian one. For example, if $\Psi_n(x_1, \ldots, x_n, t)$ is the *n*-particle wavefunction (the projection of the Fock state onto the *n*-sector), the *i*th particle's velocity is given by the Bohmian guidance formula

$$\frac{d\mathbf{x}_i}{dt}(t) = \frac{\hbar}{m_i} \Im\left(\frac{\nabla_{\mathbf{x}_i}\Psi_n}{\Psi_n}\right)(x_1, \ldots, x_n, t), \tag{1}$$

for spinless particles, in appropriate units. This "free motion" part of the dynamics corresponds to the free Hamiltonian H_0 of the QFT, and is completely analogous to the standard pilot-wave theory (or Bohm–Dirac theory for relativistic fermions) where the particle number is fixed. It is useful to recall briefly how (1) arises from the quantum probability current. On a fixed *n*-particle sector, the wavefunction Ψ_n obeys the Schrödinger equation

$$i\hbar \frac{\partial \Psi_n}{\partial t} = H_0 \Psi_n, \tag{2}$$

with H_0 a standard *n*-particle Hamiltonian. From this one obtains the familiar continuity equation

$$\frac{\partial |\Psi_n|^2}{\partial t} + \sum_{i=1}^{n} \nabla_{\mathbf{x}_i} \cdot \mathrm{j}_i = 0, \qquad \mathbf{j}_i = \frac{\hbar}{m_i} \Im\big(\Psi_n^* \nabla_{\mathrm{x}_i} \Psi_n\big), \tag{3}$$

so the quantum probability density $|\Psi_n|^2$ is transported by the currents $\mathbf{j}_i$ in configuration space. Let $\rho_Q(q, t)$ denote the distribution of actual configurations *Q*(*t*) on configuration space. If these configurations move along trajectories with velocities $\mathbf{v}_i(q, t)$, then ρ_Q satisfies the continuity equation

$$\frac{\partial \rho_Q}{\partial t} + \sum_{i=1}^{n} \nabla_{\mathbf{x}_i} \cdot \big(\rho_Q \, \mathbf{v}_i\big) \;=\; 0. \tag{4}$$

Equivariance is the requirement that, if $\rho_Q(\cdot, t_0) = |\Psi_n(\cdot, t_0)|^2$ for some initial time t_0, then $\rho_Q(\cdot, t) = |\Psi_n(\cdot, t)|^2$ for all later times. Comparing (3) and (4), this is achieved if $\rho_Q \, \mathbf{v}_i = \mathbf{j}_i$, and hence, whenever $\Psi_n(q, t) \neq 0$,

$$\mathbf{v}_i(q, t) \;=\; \frac{\mathbf{j}_i(q, t)}{\rho_Q(q, t)}.$$

Imposing the quantum equilibrium condition $\rho_Q = |\Psi_n|^2$ then yields

$$\mathbf{v}_i(q, t) \;=\; \frac{\mathbf{j}_i(q, t)}{|\Psi_n(q, t)|^2} \;=\; \frac{\hbar}{m_i} \, \Im\!\left(\frac{\nabla_{\mathbf{x}_i} \Psi_n}{\Psi_n}\right)(q, t),$$

which is precisely the guidance law (1). Thus the Bohmian velocities are obtained by taking the quantum probability current density and interpreting it as the flow of an actual configuration distributed according to $|\Psi_n|^2$.

2.2 Stochastic Jumps Between Sectors

Crucially, QFT contains interaction terms in the Hamiltonian that can change particle number [42]. For example, a field interaction can create a particle–antiparticle pair, and an atom can emit or absorb a photon. In the Bell-type Bohmian model, these are accounted for by allowing random jumps in the configuration $Q(t)$ between sectors of different n [11]. At random times, the configuration may "jump" from an n-particle configuration q to a new configuration q' that lies in the $n+1$ sector (for particle creation) or $n-1$ sector (for particle annihilation). Between such jumps, $Q(t)$ evolves smoothly according to the guiding equation within a single sector. Each jump corresponds to a world-line beginning or ending, while the other particle world-lines continue through the event — albeit perhaps with an abrupt change in velocity due to a change in wavefunction from Ψ_n either to Ψ_{n+1} or Ψ_{n-1} that is concomitant on the jump between sectors.

Mathematically, the jump process is defined in BQFT so as to preserve the $|\Psi|^2$ law (equivariance) on the full configuration space [9]. This imposes specific requirements on the jump rates – the probabilities per unit time for transitions between configurations. A minimal and natural choice for these rates was given by Dürr et al. [10]. In the continuum formulation,[1] a convenient expression is [10, p. 4145]:[2]

$$\sigma_\Psi(q \to q') = \frac{2}{\hbar} \, \Im \left[\frac{\Psi^*(q') \, \langle q'|H_I|q\rangle \, \Psi(q)}{\Psi^*(q)\Psi(q)} \right]_+, \tag{5}$$

[1] (Configuration) space is treated here as a continuous manifold, not a lattice as in [3], and hence transition rates are described by rate densities rather than discrete probabilities.

[2] For a more detailed derivation of this formula for the jump rate, see [12, pp. 3–10].

where $[x]_+ = \max\{0, x\}$ denotes the positive part. It is useful to sketch how (5) follows from the requirement of equivariance. Consider a pure jump Markov process on configuration space with rate density $\sigma_\Psi(q \to q')$, so that the configuration density $\rho_Q(q,t)$ evolves according to the master equation

$$\frac{\partial \rho_Q(q,t)}{\partial t} = \int dq' \Big[\sigma_\Psi(q' \to q)\, \rho_Q(q',t) - \sigma_\Psi(q \to q')\, \rho_Q(q,t)\Big]. \quad (6)$$

Equivariance requires that this evolution coincide with the Schrödinger evolution of $\rho_Q(q,t) = |\Psi_t(q)|^2$ generated by the interaction Hamiltonian H_I. For an interaction given by an integral operator

$$(H_I\Psi)(q) = \int dq'\, \langle q|H_I|q'\rangle\, \Psi(q'),$$

the Schrödinger equation implies

$$\frac{\partial |\Psi_t(q)|^2}{\partial t} = \frac{2}{\hbar} \int dq'\, \Im\Big[\Psi_t^*(q)\, \langle q|H_I|q'\rangle\, \Psi_t(q')\Big]. \quad (7)$$

The requirement that (6) with $\rho_Q = |\Psi_t|^2$ should reproduce (7) for all q and t leads, by identification of the integrands, to the constraint

$$\sigma_\Psi(q' \to q)\, |\Psi_t(q')|^2 - \sigma_\Psi(q \to q')\, |\Psi_t(q)|^2 = \frac{2}{\hbar}\, \Im\Big[\Psi_t^*(q)\, \langle q|H_I|q'\rangle\, \Psi_t(q')\Big]. \quad (8)$$

The right-hand side is antisymmetric under $q \leftrightarrow q'$. A natural choice of rates is obtained by requiring that, for every unordered pair $\{q, q'\}$, at most *one* of the two transitions $q \to q'$ or $q' \to q$ is allowed. This is achieved by setting

$$\sigma_\Psi(q \to q')\, |\Psi_t(q)|^2 = \frac{2}{\hbar}\, \Im\Big[\Psi_t^*(q')\, \langle q'|H_I|q\rangle\, \Psi_t(q)\Big]_+$$

and dividing by $|\Psi_t(q)|^2$. This formula (a continuous analogue of John Bell's lattice jump rates [3]) defines an equivariant Markov process: the factor in the numerator involves the matrix element of the interaction Hamiltonian H_I between the n-particle and m-particle sectors (with $m = n \pm 1, n \pm 2$ depending on the interaction) weighted by the appropriate wavefunction amplitudes. The denominator $\Psi^*(q)\Psi(q) = |\Psi(q)|^2$ normalises the rate by the Born-rule probability density at the current configuration. The $\Im(\cdot)$ (imaginary part) picks out the part of the transition amplitude that contributes to probability flow (ensuring detailed balance with the continuum Schrödinger evolution), and the positive part $\max\{0, \cdot\}$ guarantees non-negativity of rates. When integrated over an infinitesimal volume dq' in the target sector, $\sigma_\Psi(q \to dq')$ gives the probability for the jump into that volume per unit time.

As Dürr et al. prove, this choice of jump law, together with the continuous guiding in between, ensures that if the initial distribution of $Q(0)$ is $|\Psi_0|^2$, then at any later time the distribution of $Q(t)$ remains $|\Psi(t)|^2$, satisfying equivariance [10]. In other words, the law of motion is constructed so that it "mirrors"' the unitary evolution of the state vector at the level of particle configuration probabilities.

For example, suppose our QFT describes a species of fermion X and a boson field Y (e.g. electrons and photons). In a simple emission/absorption model, there are only two possible types of jumps: a single Y-particle appearing (emission) or disappearing (absorption), changing the particle number by ± 1.[3] In configuration space terms, H_I connects the n-particle sector to the $(n+1)$-particle sector (emission of a Y) or the $(n-1)$-particle sector (absorption of a Y).

The rate for an emission jump from configuration q (with no Y at some location) to a configuration q' that includes a new Y particle at position y is governed by the matrix element $\langle q'|H_I|q\rangle$, which might be proportional to the wavefunction amplitude for having a Y at y (and the X in the same configuration as in q). The guiding wavefunction ensures that if the quantum state has, say, a high amplitude for "X at x_0 emits a Y and goes to state $|q'\rangle$," then the Bohmian particle configuration will jump to that q' with corresponding probability.

Similarly, absorption corresponds to a jump from q (with a Y present) to q' (with one fewer particle), at a rate influenced by the amplitude for the Y to be absorbed into the field (e.g. into the kinetic energy of an X). Other processes like pair creation (a jump from n to $n+2$ particles, e.g. an electron–positron pair) or pair annihilation (n to $n-2$) can be incorporated if the interaction allows, though for simplicity many models consider only single-particle changes at a time. In all cases, the other particle coordinates remain the same across the jump, so the world-lines of other particles are continuous through the event.

2.3 Summary

To summarise, the standard Bell-type BQFT provides:

1. a state vector $\Psi(t)$ obeying the usual quantum field evolution (typically a Schrödinger-picture or multi-time wave equation),
2. a particle configuration $Q(t)$ (including particle types and positions) with a variable particle number N_t (so that when the actual configuration lies in the n-particle sector one has $N_t = n$), and
3. a hybrid law of motion: $Q(t)$ follows Bohm's deterministic trajectory equations when N_t is constant, and undergoes stochastic jumps $N_t \to N_t \pm k$ at random times with probabilities determined by $\Psi(t)$.

Like the original, non-relativistic Bohmian theory, this BQFT fully specifies individual histories of particles (thus solving the "measurement problem" conceptually by positing actual outcomes), while remaining consistent with the statistical predictions of ordinary QFT. Unlike in standard BM, however, the trajectories of these particles are not fully deterministic. Nonetheless, the randomness in the jump times and outcomes is not introduced ad hoc to mimic measurement collapse, but emerges

[3] Corresponding to type (i) and (ii) in the classification of Dürr et al. in [10, p. 4145].

naturally from requiring equivariance. The resulting picture is one in which particles really are created and annihilated: their world-lines can start or terminate, and these events occur with well-defined probabilities given the state of the system.

While this standard framework is already a dramatic ontological extension of QFT, it still treats the creation and annihilation of particles as spontaneous events dictated entirely by the microscopic state (Ψ, Q). In the next section, we introduce a new hypothesis: that these creation and annihilation processes also depend on *contextual, macroscopic (or, "higher-level") conditions*. We retain the overall structure of BQFT just described, but we allow the jump rates governing stochastic jumps between sectors to be context-dependent.

3 CBQFT-1: Incorporating Top-Down Effects

In standard QFT, particle transitions are typically modelled as intrinsic features of the system's microscopic Hamiltonian. However, quantum thermodynamics and the theory of open quantum systems both indicate that microscopic processes such as emission, absorption, as well as the creation and annihilation of particles, are profoundly shaped by their macroscopic context. In this section, we consider how we might build this sensitivity into BQFT to construct our first model, CBQFT-1, while keeping the discussion heuristic. In Section 4, we will formalise this model.

3.1 Defining Context

We shall begin by defining a *macroscopic context* as any collection Λ of classical variables that influence the system's dynamics but are not themselves part of the system's quantum state. Later on, in Section 7, we will refine this definition in the course of formalising our second model, CBQFT-2.

3.1.1 Two Kinds of Macroscopic Contexts

For the purposes of this paper, two broad classes of context are considered:[4]

1. Thermodynamic environment This refers to properties of the ambient environment (e.g. temperature, entropy gradients, electromagnetic field background, radiation bath, etc.) which characterise the system's coupling to a heat bath or external fields. These thermodynamic context variables can often be treated as classical parameters (e.g. a classical electromagnetic field amplitude E, or the fixed temperature T of a reservoir).

Let's take stock of a few significant cases where the influence of thermodynamic context on quantum dynamics in modulating, constraining, or even initiating quantum events, is both experimentally and theoretically well-established. For instance,

[4] Λ need not be thermodynamic or at equilibrium. In our examples we discuss thermal contexts because they allow a transparent detailed-balance illustration.

it is well-established in the field of quantum optics that the temperature of the surrounding environment directly influences emission and absorption processes. An excited atom in a thermal field is not only subject to spontaneous emission but also to stimulated emission, with a rate enhanced by the photon occupation number $\bar{n}(\omega, T)$, which depends on the Planck distribution [36, 47].[5] It is also known that entropy gradients in non-equilibrium environments can drive irreversible dynamics and directional feedback in quantum systems [17, 23]. When coupled to quantum degrees of freedom, these gradients can bias the transition rates between states, providing a natural avenue for top-down causation via macroscopic thermodynamic constraints. In similar fashion, chemical potentials and macroscopic boundary conditions in semiconductors and condensed matter systems alter the statistical structure of quantum ensembles, particularly those obeying Fermi-Dirac or Bose-Einstein statistics [39]. As a final example, take the Schwinger effect, where electric fields enable spontaneous electron-positron pair creation from the vacuum [8, 50], modifying the vacuum structure and particle dynamics. Such field-induced effects are crucial in understanding how classical field configurations can exert top-down influence on quantum field behaviour.

2. Apparatus context This includes the presence and configuration of measurement apparatus, boundary conditions of the experiment, or macroscopic setup that "asks" a particular question of the quantum system. In orthodox quantum mechanics, it is well known that the choice of measurement context (which observable is being measured, how the apparatus is arranged) affects the outcome probabilities – quantum phenomena are contextual in this sense. Here we mean something more concrete: for example, the presence of a particle detector in a region of space could be represented in our model as a context that influences particle annihilation rates, since if a detector is ready to absorb a particle, the particle might have an *enhanced* chance to disappear from the system into the detector. It could alter the transition probabilities for particle annihilation by creating a "sink" into which particles may disappear. Likewise, boundary conditions set up by the apparatus (such as reflective or absorbing walls, etc.) can alter the mode structure of the field and hence modify creation/annihilation probabilities, akin to the way in which a cavity modifies emission via the Purcell effect [43].

At this heuristic stage, we are not specifying an outcome-indexed state-update rule (POVM instrument); rather, we are modelling the apparatus as a macroscopic sink that biases the particles' jump hazards and bends their trajectories. In the standard Bohmian picture, the measurement apparatus is part of the total quantum system. The measurement context refers to the choice of observable and the experimental setup, influencing outcome probabilities through entanglement and decoherence. Here, in contrast, we explore a description in which the apparatus (though still made of particles) is treated semi-classically as an *external* context that can influence the dynam-

[5] These temperature-dependent effects are captured in the open systems formalism, where Lindblad master equations incorporate the environment's influence upon the quantum system through dissipative terms [7].

ics of the primary system. In other words, in our contextual version of BQFT, we are giving the "apparatus context" greater ontological weight and more causal power than it receives in standard Bohmian theory: it physically shapes the dynamics of the quantum system, and yet it does so without collapsing the wave function. In the next section, we will explain how.[6]

3.1.2 Macroscopic Context Variables

In anticipation of our extension of Bohmian theory, let us introduce context variables $\Lambda(t)$ to denote a collection of parameters $\lambda_1, \lambda_2, \ldots$ describing things like temperature, classical fields, or measurement device settings. These variables are *classical* (i.e. not subject to quantum uncertainty within our model) and *external* (i.e. not determined by the system's wavefunction, though they may have their own dynamics or be set by initial conditions of the universe). The introduction of Λ is an additional element beyond the standard Bohmian ontology of "Ψ and Q". We will see that setting Λ to neutral or trivial values (e.g. $T = 0$, no apparatus present, etc) will reduce our theory back to the standard Bell-type BQFT described in Section 2. But for Λ corresponding to non-trivial context (e.g. finite temperature $T \neq 0$, presence of an apparatus, etc.), our hypothesis is that the particle creation and annihilation rates will be altered in a specific way.

Our approach embodies the philosophical idea of top-down causation, wherein higher-level conditions (like the presence of a macroscopic apparatus) can influence lower-level processes (like the motion of microscopic particles) – as in the contextual approach to quantum mechanics taken by Drossel and Ellis [16]. An intuitive way of understanding how top-down causation arises is to recognise that the way in which particles are *organised* in a system makes a difference to its causal powers [53]. Such influence need not violate any fundamental microphysical laws. Rather, top-down influences can enter through the boundary conditions or through additional terms that modulate those laws. In our case, the fundamental micro-law (the Bohmian jump rate formula) will be extended to include dependence on these context parameters. Effectively, we are elevating certain classical properties of the environment to the status of inputs for the microscopic evolution of the system.

3.2 Context-Dependent Jumps and Trajectories

To incorporate context into the Bohmian dynamics, we propose the following three modifications to the standard Bell-type QFT rules:

1. Wavefunction evolution with context The fundamental quantum state $\Psi(t)$ of the system still evolves primarily according to the QFT Hamiltonian $H = H_0 + H_I$ (free part plus interaction). In quantum electrodynamics, for example, the interaction term

[6] The terms "macroscopic context" and "apparatus context" should not be conflated with the term "measurement context" as used in discussions of quantum contextuality, denoting a set of commuting observables that can be jointly measured.

H_I couples the Dirac field to the electromagnetic field. However, we now allow the possibility that the Hamiltonian depends on context parameters, i.e.

$$H(\Lambda) = H_0 + H_I + H_C(\Lambda),$$

where $H_C(\Lambda)$ is an additional term that models coupling to an environment or classical background (the context). This H_C may be time-dependent, reflecting changing external conditions. Depending on the application, one may take $H_C(\Lambda)$ to be self-adjoint (as with ordinary external fields), or use complex terms as a compact phenomenological representation of loss channels (introducing a non-Hermitian component that must be handled with care).[7]

For example, if Λ includes a classical field, such as an external laser driving a transition or a background field providing energy, and the system contains a particle with an electric dipole moment $\mathbf{d}$, then H_C may include a term like $-\mathbf{d} \cdot \mathbf{E}_{\text{ext}}(t)$ [47]. A term like this models processes such as stimulated emission or absorption. In a context involving a particle detector, one often uses an effective absorption term of the form $H_C(\Lambda) = -i\Gamma(x)\, a^\dagger(x)a(x)$, where $\Gamma(x) \in \mathbb{R}_{\geq 0}$ represents a spatially localised absorption rate coincident with where the detector is placed, and $a^\dagger(x), a(x)$ are creation and annihilation operators [44]. Adding such a term to H leads to non-unitary evolution where probability "leaks" from the effective system to an environment.[8] If a detector becomes active at time t_0, one can model this via a time-switched interaction: $H(t) = H_0 + H_I + \theta(t - t_0)\, V_{\text{int}}(x)$, where θ is a Heaviside step function and $V_{\text{int}}(x)$ is a localised interaction potential [20].

2. Guiding equation (continuous motion) Between jumps, the particles still follow the guidance equation which depends on the total Hamiltonian (1). One might expect, then, that the context-dependent interaction $H_C(\Lambda)$ would also affect the continuous motion of the particles, since the Bohmian velocity field $v(q, t)$ is derived from the full Hamiltonian's action on the wavefunction (via the quantum current). Yet the dependence is somewhat subtle: as we see in (1), the continuous motion of the particles depends specifically on the spatial structure of Ψ through gradients like $\nabla_q \Psi/\Psi$. Consequently, if the context dependence in $H_C(\Lambda)$ amounts only to a spatially uniform scalar term (e.g. a purely time-dependent energy shift), then it does not change the *spatial* structure of Ψ relevant to the ratio $\nabla_q \Psi/\Psi$. In that case, the velocity field is unchanged.

3. Stochastic jump law with context The same context, however, can affect the *stochastic jumps* that govern particle creation and annihilation. We propose that the jump rate $\sigma_\Psi(q \to q')$ should be a function not only of the wavefunction Ψ and current configuration q, but also of the context Λ. We denote the rate in the new theory

[7] We distinguish these two possible regimes and explore their implications in Section 4.3.

[8] When we define the jump process more carefully in Section 4, only the Hermitian interaction responsible for probability currents on configuration space enters the jump-rate formula; anti-Hermitian contributions encode norm loss but do not generate jumps.

by $\sigma_{\Psi,\Lambda}(q \to q')$. One heuristic but convenient way to write this is as a context-weighted version of the standard rate (5). For instance, we can write

$$\sigma_{\Psi,\Lambda}(q \to q') \;=\; g_\Lambda(q \to q')\, \sigma_\Psi^{(0)}(q \to q'),$$

where $\sigma_\Psi^{(0)}$ is the original jump rate (with no context) and $g_\Lambda(q \to q')$ is a non-negative function (a contextual modifier) that encapsulates the influence of context Λ on that transition. In the absence of context, $g_\Lambda \equiv 1$, restoring $\sigma_{\Psi,\Lambda} = \sigma_\Psi^{(0)}$. When context is present, g_Λ might either enhance ($g > 1$) or suppress ($0 \leq g < 1$) the base rate, depending on the nature of the influence.

Alternatively—and more precisely—one can write a formal context-extended analogue of (5) by including H_C. If $H = H_0 + H_I + H_C(\Lambda)$, then we have:[9]

$$\sigma_{\Psi,\Lambda}(q \to q') = \frac{2}{\hbar}\, \Im \left[\frac{\Psi^*(q')\, \langle q'|H_I + H_C(\Lambda)|q\rangle\, \Psi(q)}{\Psi^*(q)\Psi(q)} \right]_+ \tag{9}$$

Jump rates are computed from matrix elements of the interaction Hamiltonian between different particle-number sectors, and these elements can be modulated by macroscopic parameters even when the corresponding terms are position-independent. For instance, a detector's readiness to absorb a particle, or the presence of a high-temperature environment, may enhance transition amplitudes between sectors of different particle number without affecting spatial evolution. Therefore, even when context leaves the trajectories untouched, it may alter the landscape of particle creation and annihilation by contributing additional matrix elements for transitions.

It is important to stress that $\sigma_{\Psi,\Lambda}(q \to q')$ still depends on the wavefunction Ψ in a crucial way (through the matrix elements $\langle q'|H_I|q\rangle$ etc.). We are not replacing the quantum probabilistic mechanism with an arbitrary classical one; rather, we are introducing a coupling between the wavefunction-driven dynamics and the context. In the limit of no context, the dynamics is purely wavefunction-driven (bottom-up, as in Section 2). With context, certain transitions might be amplified or suppressed, but only those transitions that are allowed and "suggested" by the wavefunction can occur at all. For example, if the wavefunction amplitude $\Psi(q')$ is zero, no context factor can induce a jump to q' (since the base rate $\sigma_\Psi^{(0)}$ is zero). Context can only modulate non-zero transition amplitudes, not create entirely new possibilities ex nihilo. In this way, the primacy of the wavefunction in guiding the process is maintained, consistent with the spirit of BM that the wavefunction never "curls up and dies" but acts as a pilot-wave (here, guiding *both* the continuous motion *and* stochastic jumps, as in Section 2).

[9] In the formalisation in Section 4, we choose to restrict $H_C(\Lambda)$ in (9) to $H_C^H(\Lambda)$ (i.e. the Hermitian part): anti-Hermitian terms produce norm loss but do not generate jumps.

3.3 Examples of Context-Dependent Jump Dynamics

CBQFT-1 is intentionally minimalist: it keeps the Bell-type primitive ontology and jump structure, while allowing macroscopic context Λ to modulate the dynamics in a controlled way. We illustrate the idea with two examples.

1. Thermally stimulated transitions For example, if Λ represents a reservoir at temperature T, the term $H_C(\Lambda)$ may effectively encode a stimulated transition amplitude between different particle-number sectors.[10] A plausible model is that the Hermitian, context-dependent interaction provides an amplitude for absorption from, or emission into, the environment. In the familiar case of a finite-temperature environment, one expects that upward transitions (creating a particle of energy E) are easier when the environment can supply energy E (thermal activation), whereas downward transitions (losing a particle of energy E) are more likely when the environment can accept energy.

In a system coupled to a thermal environment at temperature T, the transition rates should satisfy the principle of *detailed balance* [7, 57].[11] This condition ensures that, in thermal equilibrium, the probability current between any two configurations q and q' vanishes, thereby maintaining a stationary probability distribution. Formally, if $\rho(q)$ denotes the equilibrium probability of configuration q, and $\tilde{g}_\Lambda(q \to q')$ is the effective transition rate (or rate density) from q to q' under the context Λ, this condition requires:

$$\rho(q)\, \tilde{g}_\Lambda(q \to q') = \rho(q')\, \tilde{g}_\Lambda(q' \to q), \qquad \text{where } \tilde{g}_\Lambda = g_\Lambda\, \sigma_\Psi^{(0)}.$$

When the system is in thermal equilibrium, the probability distribution follows the Boltzmann form $\rho(q) \propto \exp[-E(q)/k_B T]$, where $E(q)$ is the energy of configuration q. Substituting into the detailed-balance equation yields

$$\frac{\tilde{g}_\Lambda(q \to q')}{\tilde{g}_\Lambda(q' \to q)} = \exp\left(-\frac{\Delta E}{k_B T}\right),$$

where $\Delta E = E(q') - E(q)$ is the energy gained in the transition from q to q'. This relation implies that upward transitions in energy are exponentially suppressed relative to downward transitions, with the suppression governed by the thermal energy scale $k_B T$. As such, the thermal context encoded in Λ directly modulates the relative likelihood of particle creation and annihilation events, particularly in the stochastic jump sector of contextual QFTs. At zero temperature, upward transitions are sup-

[10] In the formal model of Section 4, only the Hermitian part of the interaction Hamiltonian contributes to the Bell-type currents entering the jump rates; anti-Hermitian parts, if they are admitted, govern norm loss at the ensemble level.

[11] We can extend to non-thermal, non-equilibrium contexts simply by dropping detailed-balance constraints and treating g_Λ and/or $\langle q'|H_C(\Lambda)|q\rangle$ as determined by the (possibly stochastic) classical process $\Lambda(t)$.

pressed entirely, while at high temperatures, transitions become nearly symmetric, dominated by thermal fluctuations. Imposing this form of detailed balance provides a natural thermodynamic constraint on the context-dependent jump dynamics and helps ensure consistency with equilibrium statistical mechanics.

Within CBQFT-1, such thermodynamic constraints can be implemented either through the context-sensitive Hamiltonian term $H_C(\Lambda)$, or by encoding them in the contextual modifier g_Λ that multiplies the standard Bell rate $\sigma_\Psi^{(0)}(q \to q')$. For instance, one might include in $H_C(\Lambda)$ a coupling that reflects the environmental mode occupancy $n_{\mathrm{env}}(\omega)$, such as the thermal photon number or phonon occupation number. This factor would then modulate the effective transition amplitudes and hence the jump rates. In such cases, the context-dependent jump rates $\tilde{g}_\Lambda(q \to q')$ would exhibit the familiar thermal enhancements and suppressions: for example, factors of $(1 + n_{\mathrm{env}})$ for boson emission, and n_{env} for absorption, in accordance with known results from quantum optics and quantum statistical mechanics [7, 21, 47].

2. Apparatus context Suppose Λ represents an apparatus that is primed to register a particle of type *Y* in a region *R* (e.g. a detector covering *R*). We can model this readiness by a context-dependent enhancement of those inter-configuration couplings that implement the relevant registration transition, so that the corresponding Bell-type jump rates out of configurations with a *Y*-particle in *R* are increased. In terms of the modified Bell law, this could be achieved with a contextual factor $g_\Lambda(q \to q')$ that is, say, $\gg 1$ if the jump $q \to q'$ corresponds to a particle hitting the detector and disappearing (since the apparatus is designed to catch it), and $g_\Lambda \approx 1$ for other jumps. In this way, the mere presence of the detector biases the Bohmian jumps so that whenever the particle reaches the detector region, it will almost certainly undergo an annihilation jump (i.e. be detected). Conversely, if a context implies a closed system (no measurement, reflective walls), we might set $g_\Lambda = 0$ for any annihilation into the environment (no particle can escape), thereby enforcing particle-number conservation if appropriate.

4 CBQFT-1: Formal Model

Having discussed the basic procedure for introducing contextual parameters into BQFT in Section 3, and hence for incorporating top-down causation, we now turn to the formal structure. This section specifies the dynamics rigorously and demonstrates how the standard non-contextual theory is recovered in limiting cases. Our aim in this section is to establish internal consistency and clarify how contextual modifications enter at the level of jump rates, Hamiltonians, and particle trajectories. In Section 5, we discuss the broader implications of introducing such contextual elements into quantum theory.

4.1 Axioms of the Model

(A1) **State space.** At each time *t*, the total physical state of the system is a triple (Ψ_t, Q_t, Λ), where:

- $\Psi(t) \in \mathcal{F}$ is the wavefunction, evolving in a fixed Fock space $\mathcal{F} = \bigoplus_n \mathcal{H}_n$.
- *Q*(*t*) is the actual configuration, including particle number *N*(*t*) and positions $\mathrm{x}_1, \ldots, \mathbf{x}_{N(t)} \in Q$.
- $\Lambda = (\lambda_1, \ldots, \lambda_M)$ is a collection of classical parameters (typically real-valued fields) that are externally specified.[12]

(A2) **Wave function dynamics.** The wavefunction evolves according to a Schrödinger-type equation with context-sensitive Hamiltonian:

$$i\hbar \frac{d}{dt}\Psi(t) = [H_0 + H_I + H_C(\Lambda)]\,\Psi(t), \tag{10}$$

where H_0 is the free Hamiltonian, H_I describes standard field interactions, and $H_C(\Lambda)$ is a context-dependent term. When $H_C(\Lambda)$ is not Hermitian, write the canonical decomposition

$$H_C(\Lambda) \;=\; H_C^{\mathrm{H}}(\Lambda) \;-\; \tfrac{i}{2}\,\Gamma_\Lambda, \qquad H_C^{\mathrm{H}}(\Lambda) = H_C^{\mathrm{H}}(\Lambda)^\dagger, \;\; \Gamma_\Lambda \geq 0.$$

(A3) **Configuration guidance.** In a sector with *n* particles and no jump occurring, the configuration evolves according to:

$$\frac{d\mathrm{x}_i}{dt} = \frac{\hbar}{m_i}\Im\left[\frac{\nabla_{\mathrm{x}_i}\Psi(q,t)}{\Psi(q,t)}\right], \qquad i = 1, \ldots, n, \tag{11}$$

with $q = (\mathrm{x}_1, \ldots, \mathrm{x}_n)$ the particle configuration, as in (1).

(A4) **Stochastic particle jumps.**
The inter-configuration coupling part of the Hamiltonian in (A2) typically has off-diagonal matrix elements $\langle q|H_I + H_C^{\mathrm{H}}(\Lambda)|q'\rangle$ with $q \neq q'$ that generate inter-configuration probability currents:[13]

$$J_\Lambda(q, q'; t) := \frac{2}{\hbar}\,\Im\big(\Psi_t^*(q)\,\langle q|H_I + H_C^{\mathrm{H}}(\Lambda)|q'\rangle\,\Psi_t(q')\big).$$

[12] Only the induced operator $H_C(\Lambda)$ enters the dynamics, so complex-valued components are permitted if desired.

[13] For a discussion of why only Hermitian terms enter the inter-configuration probability currents and jump rates in (A4), see Remarks 4.3 and 4.4.

Particle jump rates are fixed so that the net gain–loss current associated with the configuration matches $J_\Lambda(q, q'; t)$ (the balance condition).[14] When in configuration q (in some sector with n particles), the rate for a jump to another configuration $q' \neq q$ (possibly within the same n sector, or a sector with $n \pm 1$) is:

$$\sigma_{\Psi,\Lambda}(q \to q') := \frac{2}{\hbar}\Im\left[\frac{\Psi^*(q')\ \left(\langle q'|H_I|q\rangle + \langle q'|\, H_C^{\mathrm{H}}(\Lambda)\,|q\rangle\right)\ \Psi(q)}{\Psi^*(q)\Psi(q)}\right]_+ . \tag{12}$$

The total escape rate is:

$$\tau(q) := \sum_{q' \neq q} \sigma_{\Psi,\Lambda}(q \to q'). \tag{13}$$

The probability of jumping to a configuration in dq' is:

$$\mathbb{P}(q \to q' \in dq') = \frac{\sigma_{\Psi,\Lambda}(q \to q')}{\tau(q)}\, dq'. \tag{14}$$

4.2 Remarks and Assumptions

Remark 4.1 *(Relation to standard Bell-type QFT) If $H_C(\Lambda) = 0$, the above axioms reduce to the usual Bell-type Bohmian QFT on a fixed Fock space: a single global wavefunction evolving under $H_0 + H_I$, a configuration Q_t guided by the Bohmian law, and stochastic jumps with rates determined solely by H_I. The role of the context term $H_C(\Lambda)$ in CBQFT-1 is to add a controlled, externally specified modulation of both the continuous and jump dynamics, while leaving the underlying Hilbert-space structure unchanged.*

Remark 4.2 *(External time dependence of the context) In CBQFT-1 the context Λ is not a dynamical variable of the theory: it does not respond to Q_t and is not subject to its own equations of motion. However, it may be externally prescribed as a function of time $\Lambda(t)$ – representing, for example, a detector being switched on or off according to a classical control protocol. All appearances of $H_C(\Lambda)$ in the axioms may thus be read as $H_C(\Lambda(t))$ if an explicit time dependence is required. Distinct values of Λ that induce the same operator $H_C(\Lambda)$ (and hence the same microscopic law) are physically equivalent as far as this model is concerned; Λ is best understood here as a coarse-grained macroscopic descriptor of experimental conditions, which is specified exogenously rather than inferred from the jump rates.*

[14] Equivalently, for each $q \neq q'$ the rates are chosen so that $\sigma_{\Psi,\Lambda}(q' \to q)\,|\Psi_t(q')|^2 - \sigma_{\Psi,\Lambda}(q \to q')\,|\Psi_t(q)|^2 = J_\Lambda(q, q'; t)$. A standard minimal choice is given in (12). Note that global norm loss in Ψ does not, by itself, affect the jump hazards.

Remark 4.3 *(On the status of the Bell–type jump rates) The Bell–type rates $\sigma_{\Psi,\Lambda}(q \to q')$ in (12) are not free phenomenological parameters. Following the construction in Section 2 (see also [10, 12]), we define them by the minimal prescription that makes the gain–loss term in the Kolmogorov equation for ρ_Q reproduce the off–diagonal quantum probability currents generated by the couplings $H_I + H_C^{\mathrm{H}}(\Lambda)$. In other words, the balance condition (8), with the substitution $H_I \to H_I + H_C^{\mathrm{H}}(\Lambda)$, is part of the definition of CBQFT-1: only drift–jump processes whose rates satisfy it are admitted.*[15]

Remark 4.4 *(Absorbers and the "no double-counting" convention) If $H_C(\Lambda)$ includes an anti-Hermitian part $-\frac{i}{2}\Gamma_\Lambda$, we do not feed that part into (12) in this version of CBQFT-1; only $H_I + H_C^{\mathrm{H}}(\Lambda)$ enters the current-based jump construction. The effect of $-\frac{i}{2}\Gamma_\Lambda$ is then felt as a loss of norm which can be accounted for at the ensemble level by a standard GKLS description if desired (see Section 4.3, below). This avoids double-counting absorption in both the Hamiltonian currents and the jump hazards. Although the jump rate (12) depends on $\Psi(t)$, it depends only on ratios of quadratic expressions in $\Psi(t)$ and is therefore invariant under global rescalings of the wave function. To put it another way: H_C^{H} is "context as interaction" (redistributes probability and drives jumps); Γ_Λ is "context as leak" (removes probability).*

Remark 4.5 *(Jump channels beyond creation/annihilation) For simplicity we have written A4 in the familiar Bell-type QFT setting where the dominant discrete transitions change particle number (cf. Section 2.2). However, nothing in the balance-condition construction requires this: one may equally allow jump channels $q \to q'$ within a fixed n-particle sector (e.g. effective internal transitions or apparatus-pointer channels included in the configuration variable), provided the corresponding Hermitian couplings generate the associated off-diagonal probability currents.*

Assumption 4.1 (Quantum equilibrium) At some initial time t_0, the configuration Q_{t_0} is distributed according to the Born rule:

$$\mathbb{P}(Q_{t_0} \in dq) = |\Psi_{t_0}(q)|^2 \, dq.$$

If the dynamics preserve the Born distribution, then the law of Q_t matches $|\Psi_t|^2$ for all subsequent times.

4.3 Bohmian Properties

In CBQFT-1, the evolution of the wavefunction $\Psi(t)$ is governed by a context-sensitive Hamiltonian of the form $H = H_0 + H_I + H_C(\Lambda)$. The wavefunction Ψ_t always evolves deterministically according to the Schrödinger-type equation (10), while the

[15] We do not start from an arbitrary GKLS generator and reverse-engineer jump rates for it. Likewise, we treat a non-Hermitian term $-\frac{i}{2}\Gamma_\Lambda$ as admissible only insofar as it admits a standard GKLS completion $\Gamma_\Lambda = \sum_\alpha C_\alpha^\dagger C_\alpha$. Purely attenuative effective generators with no such channel structure are outside the intended scope of CBQFT-1.

actual configuration Q_t is a stochastic jump-drift process which is guided by Ψ_t (11) but subject to random configuration jumps (12).

At this stage there is a modelling choice. If one *restricts* to the *Hermitian regime* $H_C(\Lambda) = H_C(\Lambda)^\dagger$, then the full generator is self-adjoint, $\|\Psi_t\|$ is conserved, and the usual Bohmian reasoning about norm preservation and equivariance goes through essentially as in the standard Bell–type Bohmian QFT of Section 2. However, many practically important open-system situations—absorption by detectors, effective loss channels, or decoherence models written with complex potentials—are often modelled by allowing an anti-Hermitian component within $H_C(\Lambda)$. Admitting such terms in CBQFT-1 increases modelling flexibility, but it also changes what can be asserted *pathwise*: the guiding state Ψ_t may lose norm and strict equivariance with $|\Psi_t|^2$ need not hold in general. This is why, when non-Hermitian context terms are permitted, it becomes useful to distinguish:

(i) *Pathwise* statements, which concern a single realised history $t \mapsto Q_t$ of the configuration, guided by the (deterministic) wavefunction Ψ_t via A3 between jumps and undergoing discrete jumps according to A4.
(ii) *Ensemble* statements, which concern statistics over many runs of the experiment, all with the same initial wavefunction Ψ_{t_0} but possibly different initial configurations and jump histories.

We denote the law of the beables by

$$\rho_Q(q,t)\,dq := \mathbb{P}(Q_t \in dq).$$

Independently of the guiding state Ψ_t used for Bohmian guidance (A3), it is often convenient when $H_C(\Lambda)$ contains non-Hermitian components to summarise the *ensemble* evolution by an effective density matrix ρ_t on $\mathcal{F}$, i.e.

$$\rho_t := \mathbb{E}(|\Phi_t\rangle\langle\Phi_t|)\,. \tag{15}$$

Here $|\Phi_t\rangle$ is an *auxiliary* stochastic state used only to represent the effective GKLS ensemble dynamics by an unravelling and should not be conflated with the guiding state Ψ_t appearing in Axioms A2–A4.[16] Although the guiding state Ψ_t is a pure vector along each realised history, averaging over the auxiliary stochastic trajectories $|\Phi_t\rangle$ typically produces a mixed ρ_t (i.e. $\mathrm{tr}(\rho_t^2) < (\mathrm{tr}\,\rho_t)^2$), which reflects coarse-graining over environmental/loss channels in the effective description. To connect this effective description back to the primitive ontology, we require for empirical adequacy

$$\rho_Q(q,t) \;=\; \langle q|\rho_t|q\rangle, \tag{16}$$

[16] One may choose an unravelling in which $|\Phi_t\rangle$ evolves with H_{eff} between *channel jumps* (Kraus updates) that encode the coarse-grained absorption/decoherence processes. These channel jumps are part of the effective open-system bookkeeping and are not, in general, identical to the configuration jumps of Q_t generated by the Bell–type rates from the Hermitian inter-configuration couplings.

i.e. the diagonal of the effective statistical state reproduces the configuration of the actual distribution.[17] Because ρ_t is defined by (15) (not by (16)), trace preservation of ρ_t is a substantive claim, along with ensemble equivariance, and each is discussed in Propositions 4.1 and 4.2 respectively below.[18]

4.3.1 Norm Conservation

When $H_C(\Lambda)$ is Hermitian, the full Hamiltonian is self-adjoint and hence the Schrödinger evolution conserves the norm $\|\Psi_t\|$. When $H_C(\Lambda)$ contains non-Hermitian components, however, amplitude flows from the explicitly modelled degrees of freedom into effective absorption/decoherence channels, represented in CBQFT-1 by the anti-Hermitian part $-\frac{i}{2}\Gamma_\Lambda$ of $H_C(\Lambda)$.[19] We deal with this situation by adopting an *open system regime* in which one averages over which channel events occurred, obtaining an unconditional mixed state ρ_t whose evolution is *trace preserving*.

Proposition 4.1 *(Ensemble norm conservation) Suppose $H_C(\Lambda)$ includes non-Hermitian contributions. Then:*

(a) *the norm of the guiding state Ψ_t solving (10) is in general not conserved pathwise; but*
(b) *norm conservation is recovered at the ensemble level: the effective density matrix ρ_t defined in (15) is trace preserving,*

$$\mathrm{tr}\rho_t = \mathrm{tr}\rho_{t_0} \qquad \text{for all } t \geq t_0.$$

Proof sketch Write $H_C(\Lambda) = H_C^{\mathrm{H}}(\Lambda) - \frac{i}{2}\Gamma_\Lambda$ and set $H_\Lambda := H_0 + H_I + H_C^{\mathrm{H}}(\Lambda)$. Then the Schrödinger-type equation (10) may be written with the (generally non-self-adjoint) generator $H_{\text{eff}} := H_\Lambda - \frac{i}{2}\Gamma_\Lambda$:

$$H_{\text{eff}} = H_\Lambda - \tfrac{i}{2}\Gamma_\Lambda, \qquad \Gamma_\Lambda = \sum_\alpha C_\alpha^\dagger C_\alpha, \quad \Gamma_\Lambda \geq 0. \tag{17}$$

This immediately yields (a): for the guiding state Ψ_t,

$$\frac{d}{dt}\|\Psi_t\|^2 = -\frac{1}{\hbar}\langle \Psi_t, \Gamma_\Lambda \Psi_t \rangle \leq 0,$$

[17] When one explicitly constructs $|\Phi_t\rangle$ as an unravelling of the same GKLS dynamics used to model the absorption channels, (16) is the usual statement that the diagonal of the statistical state yields the Born-rule distribution for configuration measurements.

[18] One does not *define* ρ_t via its diagonal. Indeed, many operators share the same diagonal, so (16) alone would not fix ρ_t.

[19] This norm loss is *not* implemented by additional jumps of the configuration Q_t in our CBQFT-1 drift–jump dynamics. The Bell-type configuration jumps are fixed by the Hermitian inter-configuration couplings $H_I + H_C^{\mathrm{H}}(\Lambda)$ (Axiom A4). The anti-Hermitian term is treated purely as an effective sink at the state/ensemble level; cf. Remark 4.4.

so $\|\Psi_t\|$ typically decreases. For the ensemble bookkeeping, we introduce an auxiliary pure-state trajectory $|\Phi_t\rangle$ evolving with the same H_{eff} between *channel* jumps.[20] Along the no-jump segments one has the analogous identity

$$\frac{d}{dt}\|\Phi_t\|^2 = -\frac{1}{\hbar}\langle \Phi_t, \Gamma_\Lambda \Phi_t\rangle,$$

so the non-Hermitian part again acts as a sink. Standard unravelling theory then restores this lost norm *at the ensemble level* by supplementing the no-jump evolution with stochastic channel jumps generated by the operators C_α [7]. The resulting ensemble density matrix $\rho_t := \mathbb{E}(|\Phi_t\rangle\langle\Phi_t|)$ obeys the GKLS master equation

$$\dot{\rho}_t = -\frac{i}{\hbar}[H_\Lambda, \rho_t] + \sum_\alpha \left(C_\alpha \rho_t C_\alpha^\dagger - \tfrac{1}{2}\{C_\alpha^\dagger C_\alpha, \rho_t\}\right), \tag{18}$$

which is trace preserving; hence $\mathrm{tr}\rho_t = \mathrm{tr}\rho_{t_0}$ for all $t \geq t_0$, proving (b). In CBQFT-1 we use ρ_t as an effective ensemble state whose diagonal reproduces the configuration distribution $\rho_Q(q,t)$. $\square$

Remark 4.6 *(Significance) In CBQFT-1, when $H_C(\Lambda)$ is non-Hermitian, norm conservation is thus an ensemble property of the effective density matrix ρ_t rather than of the single wavefunction Ψ_t. At the level of Ψ_t, norm loss represents probability flowing into absorption channels; at the ensemble level, the GKLS dynamics ensures that this probability is properly accounted for and the total trace of ρ_t is preserved.*

Remark 4.7 *(Norm loss and guidance) Context can bend Bohmian trajectories by changing the evolution of Ψ_t through $H_C(\Lambda)$. A loss of norm in CBQFT-1, however, does not as such interfere with the Bohmian guidance of the configuration Q_t, since the velocity field in (11) depends only on the ratio $\nabla\Psi/\Psi$ and is therefore invariant under arbitrary time-dependent rescalings $\Psi \mapsto c(t)\Psi$. In particular, any purely global decay factor in Ψ_t due to non-Hermitian terms affects the overall normalisation but not the trajectories. Trajectory-bending is therefore tied to those parts of $H_C(\Lambda)$ that reshape the phase structure of Ψ_t across configuration space. (See online Appendix.)*

4.3.2 Equivariance of the Configuration

Equivariance in the usual Bohmian sense means that if the initial distribution of Q_t is $\rho_Q(q,t_0) = |\Psi_{t_0}(q)|^2$, then for all later times $\rho_Q(q,t) = |\Psi_t(q)|^2$. When $H_C(\Lambda)$ is Hermitian, Bohmian equivariance holds straightforwardly. When $H_C(\Lambda)$ is non-Hermitian, however, $|\Psi_t|^2$, we have to adopt a weaker version.

[20] Channel jumps are the stochastic Kraus updates generated by the operators C_α in an unravelling of the GKLS equation; they redistribute ensemble weight so that the average state ρ_t is trace preserving. They are not the same as the configuration jumps of Q_t.

Proposition 4.2 *(Equivariance of Q_t) Let $H_C(\Lambda)$ be possibly non-Hermitian, and let the configuration process evolve by Bohmian guidance between jumps and by the Bell-type jump rates of Axiom A4. If these jump rates satisfy the Bell-type gain–loss balance condition matching the off-diagonal quantum probability currents, then:*

(a) *Ensemble equivariance holds at the level of the density matrix: if the initial configuration is distributed as $|\Psi_{t_0}(q)|^2$, then*

$$\rho_Q(q,t) = \langle q|\rho_t|q\rangle \text{ for all } t \geq t_0,$$

where ρ_t is the GKLS ensemble state of the system (18).

(b) *Only in the Hermitian case, $H_C(\Lambda) = H_C(\Lambda)^\dagger$, however, does one recover full Bohmian equivariance:*

$$\rho_Q(q,t) = |\Psi_t(q)|^2 \text{ for all } t \geq t_0,$$

where $\Psi_t(q)$ is the pure-state wave function.

Proof sketch Let us begin by considering the wave function.[21] Fix a particle-number sector and a time interval on which the context Λ is constant. Within this sector, A3 implies that the Bohmian motion of the particles transports $|\Psi_t(q)|^2$ according to the usual continuity equation

$$\partial_t|\Psi_t(q)|^2 = -\nabla_q \cdot J(q,t),$$

where *J*(*q*, *t*) is the standard Bohmian current. This term represents probability flow *within* the sector arising from particle motion. To incorporate the effects of the full Hamiltonian, which admits interaction and a contextual influence, it is convenient to write the (possibly non-Hermitian) generator as

$$H_{\text{eff}} = H_0 + H_I + H_C^{\text{H}}(\Lambda) - \tfrac{i}{2}\Gamma_\Lambda,$$

as before (17), and work in the configuration basis. A standard calculation starting from $i\hbar\partial_t\Psi_t = H_{\text{eff}}\Psi_t$ and its adjoint gives

$$\partial_t|\Psi_t(q)|^2 = \frac{2}{\hbar}\sum_{q'} \Im\Big(\Psi_t^*(q) H_{qq'}^{\text{eff}} \Psi_t(q')\Big),$$

where $H_{qq'}^{\text{eff}}$ can be decomposed into a part coming from H_0 that reproduces the drift term $-\nabla_q \cdot J(q,t)$, as above, a Hermitian part coming from the couplings $H_I + H_C^{\text{H}}(\Lambda)$ that produce currents between configurations and enter as $\sum_{q'\neq q} J_{\text{H}}(q,q';t)$, and an anti-Hermitian part $-\frac{i}{2}\Gamma_\Lambda$ which contributes a local sink term $-J_\Gamma(q,t)$ with $J_\Gamma(q,t) \geq 0$, encoding loss of norm due to absorption. Hence the full quantum continuity balance can be written as

[21] Cf. the discussion in [12], Section 2.2, pp.4-6.

$$\partial_t |\Psi_t(q)|^2 \;=\; -\nabla_q \cdot J(q,t) \;+ \sum_{q' \neq q} J_{\mathrm{H}}(q,q';t) \;-\; J_\Gamma(q,t). \tag{19}$$

Let us turn now to the particles. Let $\rho_Q(q,t)\,dq := \mathbb{P}(Q_t \in dq)$. Because Q_t is a drift–jump process, involving Bohmian guidance between jumps, plus stochastic creation/annihilation jumps, ρ_Q obeys the Kolmogorov equation

$$\partial_t \rho_Q(q,t) = -\nabla_q \cdot J(q,t) + \sum_{q'} \Big[\sigma_{\Psi,\Lambda}(q' \to q)\,\rho_Q(q',t) - \sigma_{\Psi,\Lambda}(q \to q')\,\rho_Q(q,t)\Big]. \tag{20}$$

The first term is the same Bohmian drift $-\nabla_q \cdot J$, while the second term is the usual gain–loss structure for a jump process with rates $\sigma_{\Psi,\Lambda}(q \to q')$, as defined in (12). It is useful to view the right-hand side as a linear operator $\mathcal{L}_t[\mu]$ acting on densities μ, so that the Kolmogorov equation reads $\partial_t \rho_Q = \mathcal{L}_t[\rho_Q]$. The Bell-type rates $\sigma_{\Psi,\Lambda}(q \to q')$ are not arbitrary: they are defined in terms of Ψ_t and the Hermitian couplings $H_I + H_C^{\mathrm{H}}(\Lambda)$ so as to satisfy the balance condition

$$\sigma_{\Psi,\Lambda}(q' \to q)\,|\Psi_t(q')|^2 - \sigma_{\Psi,\Lambda}(q \to q')\,|\Psi_t(q)|^2 \;=\; J_{\mathrm{H}}(q,q';t),$$

for all $q \neq q'$. This is the natural context-generalisation of the Bell–type gain–loss condition derived in Section 2.2 for the standard jump rates $\sigma_\Psi(q \to q')$. This is a constraint on the rates themselves and does *not* presuppose any relation between ρ_Q and $|\Psi_t|^2$. To test whether $|\Psi_t|^2$ is transported in the same way as ρ_Q by the drift–jump dynamics, we now evaluate the Kolmogorov generator on the specific choice $\mu(q,t) := |\Psi_t(q)|^2$. Using the balance condition, the jump term in $\mathcal{L}_t[\mu]$ becomes

$$\sum_{q'} \Big[\sigma_{\Psi,\Lambda}(q' \to q)\,|\Psi_t(q')|^2 - \sigma_{\Psi,\Lambda}(q \to q')\,|\Psi_t(q)|^2\Big] = \sum_{q' \neq q} J_{\mathrm{H}}(q,q';t),$$

so that

$$\mathcal{L}_t[|\Psi_t|^2](q) = -\nabla_q \cdot J(q,t) + \sum_{q' \neq q} J_{\mathrm{H}}(q,q';t).$$

In other words, the Kolmogorov generator acting on $|\Psi_t|^2$ reproduces exactly the drift term $-\nabla_q \cdot J$ and the Hermitian inter-configuration currents $\sum_{q' \neq q} J_{\mathrm{H}}(q,q';t)$ that appear in the quantum continuity equation above. Comparing the two balances, we find

$$\partial_t |\Psi_t(q)|^2 - \mathcal{L}_t[|\Psi_t|^2](q) \;=\; -\,J_\Gamma(q,t).$$

When $H_C(\Lambda)$ is Hermitian, $\Gamma_\Lambda = 0$ and $J_\Gamma \equiv 0$, so this reduces to

$$\partial_t |\Psi_t(q)|^2 = \mathcal{L}_t[|\Psi_t|^2](q).$$

On the other hand, the Kolmogorov equation for the configuration law reads

$$\partial_t \rho_Q(q,t) = \mathcal{L}_t[\rho_Q](q).$$

Thus, in the Hermitian case, $|\Psi_t|^2$ and ρ_Q satisfy the same linear forward equation with the same generator $\mathcal{L}_t$. If the initial configuration is distributed according to the Born rule, $\rho_Q(q,t_0) = |\Psi_{t_0}(q)|^2$, it follows that $\rho_Q(q,t) = |\Psi_t(q)|^2$ for all $t \geq t_0$; the pathwise Bohmian equivariance asserted in (b).

In the non-Hermitian case, however, the sink J_Γ prevents $|\Psi_t|^2$ from satisfying the Kolmogorov equation, so ρ_Q no longer tracks $|\Psi_t|^2$ pathwise. At this point one switches to the ensemble description we discussed earlier. As in Proposition 4.1, the non-Hermitian contribution can be written as $\Gamma_\Lambda = \sum_\alpha C_\alpha^\dagger C_\alpha$, and standard results from open quantum systems show that the corresponding ensemble density matrix ρ_t obeys the GKLS equation (18), which is trace preserving. The diagonal $\langle q|\rho_t|q\rangle$ then obeys the same forward equation as $\rho_Q(q,t)$, so if they agree initially they agree for all later times—this is *ensemble equivariance*, proving (a). □

Remark 4.8 *(Pathwise vs. ensemble equivariance) In CBQFT-1, any non-Hermitian context terms cause the single wavefunction Ψ_t to lose norm in time. Since the usual Bohmian equivariance relation identifies the configuration density with $|\Psi_t|^2$, this norm loss disrupts the standard continuity argument: $|\Psi_t|^2$ no longer represents a normalised probability density, even though the law of Q_t must remain normalised. In this strict pathwise sense, $\rho_Q(q,t) = |\Psi_t(q)|^2$ is not automatic once $H_C(\Lambda)$ becomes non-Hermitian. Consistency is restored at the ensemble level: the probability lost is recovered in the GKLS channel terms, so that ρ_t is trace preserving and ρ_Q can be identified with $\langle q|\rho_t|q\rangle$. In this sense, a generalised equivariance is maintained even in the presence of non-Hermitian context terms: the configuration distribution is equivariant with respect to the ensemble state ρ_t, and reduces to the usual Bohmian equivariance $\rho_Q = |\Psi_t|^2$ whenever $H_C(\Lambda)$ is Hermitian. This loss of strict pathwise equivariance under these conditions, however, is one of the motivations for our second model, CBQFT-2, where the guiding state remains normalised and full Bohmian equivariance holds (Sections 6-8).*

4.4 Limiting Regimes

In the weak-coupling limit where $H_C = \epsilon V$ with $\epsilon \to 0$, and V is Hermitian, the jump rate expands as:

$$\sigma_{\Psi,\Lambda}(q \to q') \approx \sigma_\Psi^{(0)}(q \to q') + \epsilon \cdot \frac{2}{\hbar} \Im \left[\frac{\Psi^*(q')\langle q'|V|q\rangle \Psi(q)}{\Psi^*(q)\Psi(q)} \right]. \tag{21}$$

This recovers standard Bell-type BQFT jump rates when $\epsilon = 0$ or when $H_C(\Lambda) = 0$; in particular, a "vacuum" context corresponds to vanishing H_C and reduces CBQFT-1 to the usual non-contextual Bell-type dynamics.

4.5 Summary

CBQFT-1 introduces a minimal extension of Bell-type Bohmian quantum field theory, incorporating macroscopic context through a fixed or externally specified parameter Λ. This parameter modulates the Hamiltonian and the jump dynamics without altering the underlying Hilbert space structure. The theory allows contextual information—such as detector settings and classical sources—to influence the quantum evolution of $\Psi(t)$ and the stochastic structure of particle transitions.

4.5.1 Benefits

This approach has several benefits:

- It preserves much of the formal structure of standard BQFT, including a fixed Fock space, a single global wavefunction, and conventional particle guidance laws.
- It introduces top-down causation within quantum field theory in a controlled and analytically tractable way, by letting macroscopic context parameters Λ systematically modulate both the effective generator and the Bell-type jump structure.
- It permits weak, semi-classical modelling of measurement, absorption, or coarse-grained environmental effects via context-dependent Hamiltonian terms and jump rates.

4.5.2 Limitations

However, CBQFT-1 also has structural limitations:

- The context variable Λ is not a dynamical element of the theory—it does not evolve stochastically or respond to particle-level events.
- The theory cannot accommodate transitions between inequivalent quantum sectors or represent physical processes like symmetry-breaking or phase transitions in a structurally faithful way.
- If $H_C(\Lambda)$ is allowed to contain non-Hermitian components, then norm is not conserved at the level of the single guiding wavefunction and the usual Bohmian equivariance relation holds only in a weakened, ensemble-level form (reflecting the open-system character of that regime).
- Not every effective open-system description can be realised as a CBQFT-1 model; in particular, purely attenuative dynamics with no matching jump channels fall outside of this framework.

In short, we are disinclined to view CBQFT-1 as a candidate fundamental Bohmian QFT: in particular, the dynamics is not *closed* in the relevant sense. It specifies a law for (Ψ_t, Q_t) that is *conditional on* an externally supplied Λ (or protocol $\Lambda(t)$), and does not supply a law for Λ itself; as a result, the theory does not purport to describe the full world without additional dynamics for the macroscopic context not included within the model.

Despite these limitations, however, CBQFT-1 is conceptually and pragmatically valuable. It captures leading-order contextual effects within a fixed representational framework, and it serves as an effective approximation to the richer and more structurally complete dynamics of CBQFT-2, where the context becomes dynamical (closing the theory at the level of (Ψ_t, Q_t, Λ_t)), Hamiltonians remain self-adjoint, and full Bohmian equivariance is restored—as we shall see shortly, in Sections 6–7.

5 CBQFT-1: Implications

One of the striking implications of CBQFT is that it relaxes the strict hierarchy of scales often assumed in physics. In conventional thinking, macroscopic behaviour emerges from the collective dynamics of microscopic constituents—a bottom-up paradigm. Our first model, CBQFT-1, allows a limited reversal: macroscopic conditions (such as the presence of a heat bath or a measuring device in a particular configuration) enter the microscopic law of motion through externally supplied parameters.

At the modelling level, CBQFT-1 admits two regimes. In a *Hermitian regime*, the dynamics is fully self-adjoint. In an *effective open-system regime*, one allows an anti-Hermitian component in $H_C(\Lambda)$ as a convenient phenomenological representation of absorption, with the associated ensemble bookkeeping described by GKLS techniques. We now examine how contextual influences are introduced within CBQFT-1 and what consequences they have. Since we are disinclined to regard CBQFT-1 as a candidate fundamental BQFT, however, the discussion should be read in that spirit. A more complete account of the consequences of introducing contextual influence into Bohmian theory is explored in Section 8 for CBQFT-2.

5.1 Macro-to-Micro Information Flow in CBQFT-1

In standard Bohmian mechanics (BM), the dynamics of the system are fully determined by the initial wavefunction and particle configuration together with a fixed Hamiltonian. The evolution proceeds either deterministically (as in non-relativistic BM) or stochastically (as in Bell-type quantum field theories).

CBQFT-1 modifies this picture by allowing macroscopic context to enter the quantum dynamics through an externally specified parameter Λ. This parameter determines additional terms in the Hamiltonian—specifically, the context-sensitive component $H_C(\Lambda)$ in equation (10)—and thereby modulates the Schrödinger-type evolution of Ψ_t and, through the resulting state-dependent currents, the statistics of stochastic configuration jumps (via the Bell-type rate prescription of Section 4). In the effective open-system regime, where $H_C(\Lambda)$ may be non-Hermitian, Λ may also encode coarse-grained loss channels through its anti-Hermitian part.

For example, switching on a detector at time t_0 corresponds to changing Λ at that moment (or prescribing $\Lambda(t)$). This change, although exogenous to the model, alters $H_C(\Lambda)$ and therefore alters the subsequent microscopic drift–jump law. In this sense, information flows *from the macro to the micro* in CBQFT-1: not by introducing new forces,[22] but by selecting (via Λ) which Hamiltonian and associated drift–jump statistics govern the micro-ontology.

Importantly, however, the form of top-down influence in CBQFT-1 is *unidirectional*: Λ affects the dynamics of Ψ and Q, but is not itself modified by microscopic events. There is no feedback loop in which quantum outcomes update context. CBQFT-1 is therefore best regarded as a semi-classical regime in which context-sensitive quantum dynamics are conditioned on externally supplied macroscopic data. We address this limitation in Sections 6-7, where we introduce CBQFT-2.

5.2 Measurement and the Arrow of Time in CBQFT-1

CBQFT-1 provides a transparent way of encoding measurement-like conditions within a Bohmian QFT model. In standard BM, definite outcomes are encoded in the configuration $Q(t)$, and the effective suppression of interference is typically discussed using decoherence relative to a pragmatic system–environment split [46]. In CBQFT-1, the relevant macroscopic conditions are made explicit in the contextual parameter Λ, which modulates the effective generator through $H_C(\Lambda)$ and thereby affects both the wave dynamics and the configuration statistics.

For example, switching on a photodetector corresponds to updating Λ, which in turn modifies $H_C(\Lambda)$ so that detection becomes highly probable in the detector region. No fundamental collapse is postulated; rather, the actual configuration $Q(t)$ is driven into outcome regions by the context-modulated drift–jump dynamics. In a Hermitian modelling of the apparatus, this can be done (in principle) by enlarging the explicitly modelled degrees of freedom. In an effective regime, one may instead represent absorption by allowing $H_C(\Lambda)$ to include non-Hermitian terms, with the resulting norm loss handled at the ensemble level (Section 4).

This framework also provides a natural locus for an arrow of time. Macroscopic context parameters—such as thermal gradients or dissipative environments—embed time-asymmetric structure into the jump rates $\sigma_{\Psi,\Lambda}$. When the context Λ encodes a non-equilibrium or entropy-increasing background (e.g., a thermal bath with $T > 0$), the jump rates need not satisfy detailed balance unless Λ is artificially time-reversed. This asymmetry is not statistical but structural: it reflects a thermodynamically open system with embedded directionality, consistent with the second law of thermodynamics [7, 57]. Nonetheless, insofar as CBQFT-1 is to be treated as an effective framework, its time-asymmetry must also be understood as *effective*. CBQFT-1 itself does not yet provide a dynamically complete account of that asymmetry; this is part of the motivation for CBQFT-2.

5.3 Causality in CBQFT-1

The inclusion of context variables expands BM's descriptive resources. The physical state is no longer just (Ψ, Q) but (Ψ, Q, Λ), where Λ encodes macroscopic conditions that shape the system's dynamics. In CBQFT-1, however, Λ is not a dynamical vari-

[22] As Ellis and others have emphasised [14, 16], top-down causation need not involve new fundamental forces or violations of microphysics.

able: it is fixed or externally prescribed and unaffected by microscopic events. From a causal standpoint, the model is therefore unidirectional: context influences the micro-dynamics, but feedback from particles to context is excluded by construction.

This asymmetry raises a natural question: can macroscopic systems like detectors remain unaffected by quantum events? In reality, detectors change state upon interaction with particles. Modelling such feedback requires Λ to become dynamical, with particle events updating the contextual state. CBQFT-1 does not incorporate this feature by design. It functions as a first-order approximation—a theory of one-way contextual influence—whose limitations motivate CBQFT-2.

5.4 Non-Locality and Relativity in CBQFT-1

Nonlocality is an established feature of BM: particle velocities depend on the global configuration and wavefunction. CBQFT-1 inherits this trait. If a context variable Λ encodes, say, a detector setting at a remote location, it may affect jump rates elsewhere via the entangled wavefunction—analogous to the usual nonlocal but non-signalling structure of Bell-type Bohmian theories [4, 37]. A fully relativistic formulation would likely require additional structure (e.g. a preferred foliation or a multi-time formalism) as in existing Bohmian QFT proposals [11, 33, 34, 58]. Since CBQFT-1 is only an effective regime, we do not pursue these issues here; the CBQFT-2 framework, to which we now turn, is better suited to a locality discussion in terms of local algebras and overlap structures.

6 CBQFT-2: Incorporating Feedback

In Section 4, we introduced CBQFT-1, in which macroscopic contexts influence microscopic quantum events via context-sensitive Hamiltonians and jump rates. This provided a formal basis for modelling top-down causation in Bohmian QFT. Yet genuine contextual dynamics must also incorporate *bottom-up causation*—the capacity of microscopic events to *reshape* the macroscopic context itself. For instance, when a particle triggers a detector, it alters the apparatus's physical state and thereby reshapes the very conditions under which subsequent quantum processes unfold.

To describe this bidirectional interaction, we must allow the macroscopic context to become a *dynamical variable*, capable of being updated in response to the micro-level configuration. In this section, we present (heuristically) the key ideas behind CBQFT-2: a richer model than CBQFT-1, in which microscopic configurations and macroscopic contexts *co-evolve* within a unified causal structure.

We now redefine context—not merely as a boundary condition or external parameter—but as a *physically instantiated structure* that determines the representation of the quantum field algebra [18, 25]. This shift allows us to model different macroscopic phases—such as distinct detector states, thermal regimes, or symmetry-breaking configurations—using *unitarily inequivalent representations* of the algebra of observables [45]. A configuration-dependent update rule governs stochastic transitions between these sectors, thereby completing the causal loop between micro-events and macro-structures.

On this view, collapse and decoherence emerge from the dynamics of context transitions, rather than being imposed by fiat—arising naturally within a precise quantum-field-theoretic framework. The section concludes with illustrative examples before we formalise the model in Section 7.

6.1 Defining Context (Again)

Earlier, in Section 3, we defined "context" as any set of variables or conditions describing the surroundings or experimental setup that are not part of the quantum state of the system itself, but which can influence its dynamics. Here, we refine that notion: a context is any set of macroscopic conditions that both (i) influence the system's dynamics and (ii) determine how the quantum observables are *represented*. In this extended model, context encodes both the physical surroundings of a quantum system and the structural framework within which the quantum state is realised.

In QFTs with infinitely many degrees of freedom, physical quantities—often called observables—are first defined abstractly using a C^*-algebra or von Neumann algebra. To describe a physical system, one then selects a concrete representation of this algebra on a Hilbert space. However, unlike in non-relativistic quantum mechanics, there is no unique or universal Hilbert space. Instead, there are typically many *unitarily inequivalent representations* of the same algebra. As Laura Ruetsche has emphasised, these representations correspond to distinct macroscopic phases or thermodynamic contexts [45]. The choice of representation is not merely a technicality: it reflects concrete physical features of the large-scale environment under which a quantum system is realised.

For example, in the thermodynamic limit, a ferromagnet admits two inequivalent representations for the + and − magnetised phases. Different thermal environments (e.g., hot vs. cold) require inequivalent representations [25, 29], as do symmetry-breaking phases or shifts in boundary conditions (e.g., introducing a cavity). Physicists routinely use such macroscopic features to determine the correct representation of the field algebra that anchors the quantum state to a distinct physical context [18, 19].

To formalise this, we introduce classical context variables Λ that both label macroscopic conditions and select representations of the field algebra. Let $\mathcal{A}$ denote the field algebra of observables. For each macroscopic context Λ, we associate a representation π_Λ acting on a Hilbert space $\mathcal{H}_\Lambda$, typically obtained via the GNS construction [18]. This construction yields a Hilbert space $\mathcal{H}_\Lambda$ and reference state Ω_Λ, from which one derives operators such as $a^\dagger_\Lambda(x)$ and $a_\Lambda(x)$ appropriate to that sector. The key point is that distinct values of Λ generically correspond to physically disjoint representations—describing macroscopically different behaviours or "phases"—even though they arise from the same algebraic structure. For example, Λ might encode the magnetisation direction of a spin lattice, the temperature and chemical potential of a thermal bath, or the operational state of a detector. CBQFT-2 thus builds on the insight that the algebraic structure of a quantum system must be situated within a macroscopic framework.

Crucially, we now treat Λ as a dynamical classical variable that evolves under the influence of quantum events. The full state of the system at time *t* is then (Ψ_t, Q_t, Λ_t), where $\Psi_t \in \mathcal{H}_{\Lambda_t}$ is the quantum state in the context-dependent Hilbert space, Q_t is the particle (or field) configuration, and Λ_t labels the active context. Between *context jumps*, Λ_t is constant and Ψ_t evolves according to a Schrödinger equation with Hamiltonian H_{Λ_t}. Between *particle jumps*, Q_t evolves continuously under the usual Bohmian guidance determined by Ψ_t. At the discrete jump times (of either kind) the state (Ψ_t, Q_t, Λ_t) is updated stochastically, according to rules which will be made precise in Section 7.

6.2 Context Updates: Closing the Circle

The central innovation of CBQFT-2 is that microscopic quantum events can dynamically alter the classical context. Concretely, suppose a quantum jump of type α occurs at time *t* (e.g., a particle enters a detector). We then update the macroscopic context via a rule of the form

$$\Lambda_{t^+} = F_\alpha(\Lambda_{t^-}, Q_t),$$

where F_α is a prescribed function. Such update rules are defined at the macroscopic level and are empirically grounded—reflecting how real-world apparatuses and materials respond to microscopic inputs [15, 16].

When a designated quantum jump occurs, two things happen atomically: (i) the configuration Q_t undergoes a discrete transition (e.g., a particle is added or removed); and (ii) the context updates, $\Lambda_{t^+} = F_\alpha(\Lambda_{t^-}, Q_t)$. The wavefunction is then reinterpreted as a vector in the new Hilbert space $\mathcal{H}_{\Lambda_{t^+}}$ associated with the updated representation [29].[23] In more general scenarios, context updates need not be deterministic. Each microscopic transition α can instead be associated with a *stochastic transition kernel*, specifying the conditional probability

$$P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t),$$

where $\mathcal{J}_t$ denotes relevant additional data (such as jump history or environmental variables). We formalise this structure in Section 7.

This bidirectional causal interaction—where microscopic transitions affect context, and context in turn shapes quantum evolution—constitutes a genuine feedback loop. The state (Ψ_t, Q_t, Λ_t) evolves as a coupled system: the wavefunction and configuration respond to Λ_t, and certain quantum events probabilistically update Λ_t itself.

[23] Equivalently, the jump induces a transition $\omega \to \omega'$ between states on the field algebra, with $\pi_{\Lambda_{t^+}}$ chosen such that ω' lies in its folium.

6.3 Examples of Context Updates

Two simple examples below illustrate how these ideas apply in practice.[24]

(i) A saturating detector Consider a photon detector with macroscopic pointer variable $\Lambda \in \{0, 1\}$, indicating whether it is in the "ready" ($\Lambda = 0$) or "fired" ($\Lambda = 1$) state. While $\Lambda = 0$, incoming photons see a Hamiltonian H_0 and jump structure appropriate to a ready detector. When a photon enters the detector region, a quantum jump of type $\alpha = \text{click}$ occurs, and we set

$$\Lambda_{t^+} = F_{\text{click}}(\Lambda_{t^-}, Q_t) = 1.$$

From that time onward, the field is described in the "fired" representation π_1 with Hamiltonian H_1, modelling a saturated detector that no longer responds in the same way to subsequent particles. In CBQFT-1, such a change in detector behaviour would have to be imposed externally by hand as a time-dependent change of $\Lambda(t)$; in CBQFT-2, the flip $\Lambda : 0 \to 1$ is *driven* by the actual Bohmian configuration.

(ii) A ferromagnetic phase flip Take a ferromagnet below its critical temperature T_c. Let $\Lambda \in \{+, -\}$ encode whether the macroscopic magnetisation is up or down. In the thermodynamic limit, the + and − phases are described by inequivalent GNS representations π_+ and π_- of the same field algebra [27, 45]. Suppose that, starting in the + phase, a rare fluctuation of the microscopic configuration Q_t nucleates a domain that grows and eventually inverts the macroscopic magnetisation. At that point the context variable flips,

$$\Lambda_{t^+} = F_{\text{flip}}(\Lambda_{t^-}, Q_t) = -,$$

and the field is henceforth represented on $\mathcal{H}_-$ with Hamiltonian H_-. The change in context is both dynamically triggered by the Bohmian microstate and dynamically efficacious: the new phase has different excitations, correlation structure, and effective couplings, which feed back into the subsequent evolution of Ψ_t and Q_t.

7 CBQFT-2: Formal Model

We now formulate CBQFT-2 by six axioms (B1–B6) that fix the state triple (Ψ_t, Q_t, Λ_t) and its context-indexed representations, unitary evolution within a fixed context, Bohmian guidance, stochastic particle jumps, the jump–Markov dynamics of the context, and an isometric rule for re-expressing Ψ_t across context changes. From these axioms we derive: (i) global continuity and guidance in the presence of context boundaries; (ii) norm conservation and equivariance; (iii) decoherence; and (iv) a thermodynamic arrow for the context switching process. We conclude with

[24] A more detailed example is provided in an online Appendix.

three limiting regimes and the relation between CBQFT-2 and hybrid classical–quantum models.

7.1 Axioms of the Model

(B1) **State space.** At each time *t*, the total physical state of the system is a triple (Ψ_t, Q_t, Λ_t), where:

- $\Lambda_t \in \mathcal{C}$ is a (piecewise constant, right-continuous) stochastic process (the "context") taking values in a discrete or continuous set of macroscopic context labels and obeying a law given in B5.[25]
- Q_t is a configuration variable (particle positions or field values).
- Ψ_t is the guiding state vector (specified below).

The elements of $\mathcal{C}$ label coarse-grained macroscopic contexts, such as detector settings, boundary conditions, and thermodynamic regimes. Concretely, for each $\Lambda \in \mathcal{C}$ there is a Hilbert space $\mathcal{H}_\Lambda$ and a representation $\pi_\Lambda : \mathcal{A} \to \mathcal{B}(\mathcal{H}_\Lambda)$ of the field algebra $\mathcal{A}$. The guiding state satisfies $\Psi_t \in \mathcal{H}_{\Lambda_t}$ and determines an algebraic state ω_t on $\mathcal{A}$ by

$$\omega_t(A) := \langle \Psi_t, \pi_{\Lambda_t}(A)\Psi_t \rangle, \qquad A \in \mathcal{A}.$$

(B2) **Intra-context unitary evolution.** On each open interval (t_k, t_{k+1}) in which the context is fixed, $\Lambda_t \equiv \lambda_k$, the state vector $\Psi_t \in \mathcal{H}_{\lambda_k}$ satisfies the Schrödinger equation

$$i\hbar \frac{\partial \Psi_t}{\partial t} = H_{\lambda_k}(t)\, \Psi_t, \qquad t \in (t_k, t_{k+1}),$$

where $H_{\lambda_k}(t)$ is a (possibly explicitly time–dependent) self-adjoint Hamiltonian on $\mathcal{H}_{\lambda_k}$ and may include, in addition to the "bare" field Hamiltonian, context-dependent terms representing mean fields, boundary conditions, detector couplings, etc. The corresponding propagator is unitary on $\mathcal{H}_{\lambda_k}$ for $t \in (t_k, t_{k+1})$. At the jump times t_k, however, when the context changes (B5), the state vector is updated according to B6.

(B3) **Configuration guidance.** The configuration Q_t evolves according to a context-sensitive velocity field derived from the single, global state vector $\Psi_t \in \mathcal{H}_{\Lambda_t}$. We define local densities and currents as expectations of local operator densities in the state ω_t:

$$\rho(x,t) := \langle \Psi_t, \hat{\rho}_{\Lambda_t}(x)\Psi_t \rangle, \qquad j(x,t) := \langle \Psi_t, \hat{j}_{\Lambda_t}(x)\Psi_t \rangle,$$

[25] $\mathcal{C}$ denotes a set of classical context labels, not the complex numbers $\mathbb{C}$. For simplicity, we take $\mathcal{C}$ to be finite or countable in this formalisation, corresponding to a coarse-graining into macroscopic regimes, though a continuous generalisation should be possible.

where $\hat{\rho}_{\Lambda_t}(x)$ and $\hat{j}_{\Lambda_t}(x)$ are the usual number and probability-current densities (or their appropriate QFT analogues) in the representation π_{Λ_t}. The local velocity field is then

$$v(x,t) := \frac{j(x,t)}{\rho(x,t)}, \qquad \text{where } \rho(x,t) > 0,$$

and Q_t follows the associated Bohmian trajectories. On regions of constant context, the continuity equation

$$\partial_t \rho(x,t) + \nabla \cdot j(x,t) = 0$$

holds on the domain where the operators are defined.

(B4) **Stochastic particle jumps.** For processes such as particle creation and annihilation, the configuration Q_t is a piecewise deterministic Markov process on configuration space. Fix a context value $\Lambda_t = \lambda$ and the corresponding representation π_λ on $\mathcal{H}_\lambda$, with Fock-like decomposition

$$\mathcal{H}_\lambda = \bigoplus_{n=0}^{\infty} \mathcal{H}_{\lambda,n}$$

into n-particle sectors. On any time interval on which Λ_t is constant, the dynamics of (Ψ_t, Q_t) is as follows:

- Between particle–jump times, Ψ_t evolves unitarily according to B2 with Hamiltonian H_λ, and Q_t follows the guidance law of B3.
- At a particle–jump time, Q_t undergoes an instantaneous transition $q \to q'$ with Bell–type rate

$$\sigma_{\Psi,\lambda}(q \to q') := \frac{2}{\hbar} \Im \left[\frac{\Psi_t^*(q') \, \langle q' | H_\lambda^{\text{int}} | q \rangle \, \Psi_t(q)}{\Psi_t^*(q) \Psi_t(q)} \right]_+ ,$$

where H_λ^{int} denotes the part of H_λ that couples different particle-number sectors $\mathcal{H}_{\lambda,n} \leftrightarrow \mathcal{H}_{\lambda,m}$ with $n \neq m$ (e.g. creation/annihilation terms), computed in the active representation π_λ selected by the current context.

The total escape rate from configuration q is

$$\tau_\lambda(q) := \sum_{q' \neq q} \sigma_{\Psi,\lambda}(q \to q'),$$

and the conditional distribution over targets is

$$\sigma_{\Psi,\lambda}(q \to q')/\tau_\lambda(q),$$

exactly as in A4, but with all matrix elements computed in the context-dependent representation π_λ. The wave function Ψ_t remains continuous across particle jumps; only the configuration Q_t changes.

(B5) **Context dynamics.** The context variable Λ_t is modelled as a continuous-time jump Markov process on $\mathcal{C}$, with transition probability kernel

$$P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t), \tag{22}$$

where $\mathcal{J}_t$ includes all additional data that may influence the context transition (e.g. particle currents, detector settings, or histories). The instantaneous transition rate from Λ_i to Λ_f is defined by

$$W_{\Lambda_i \to \Lambda_f}(Q_t, \mathcal{J}_t) := \lim_{\Delta t \to 0} \frac{1}{\Delta t} P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t), \qquad (\Lambda_f \neq \Lambda_i), \tag{23}$$

so that for small Δt the transition probability satisfies

$$P(\Lambda_f, t + \Delta t \mid \Lambda_i, t) = W_{\Lambda_i \to \Lambda_f}(Q_t, \mathcal{J}_t)\, \Delta t \; + \; o(\Delta t). \tag{24}$$

(B6) **Isometric re-expression.** If Λ jumps from Λ_i to Λ_f, there exist an overlap algebra $\mathcal{A}_0 \subset \mathcal{A}$ and a (partial) isometry

$$U_{\Lambda_i \to \Lambda_f} : \mathcal{H}_{\Lambda_i} \to \widetilde{\mathcal{H}}_{\Lambda_f}$$

such that for all $A \in \mathcal{A}_0$ and all Ψ in the initial subspace $\overline{\pi_{\Lambda_i}(\mathcal{A}_0)\Psi_{t^-}} \subset \mathcal{H}_{\Lambda_i}$,[26]

$$\langle \Psi, \pi_{\Lambda_i}(A)\Psi \rangle = \left\langle U_{\Lambda_i \to \Lambda_f} \Psi,\, \widetilde{\pi}_{\Lambda_f}(A)\, U_{\Lambda_i \to \Lambda_f} \Psi \right\rangle \quad \forall\, A \in \mathcal{A}_0.$$

Upon a context jump $\Lambda_i \to \Lambda_f$ at time *t*, the guiding state is updated:

$$\Psi_{t^+} = U_{\Lambda_i \to \Lambda_f} \Psi_{t^-}.$$

There are two kinds of context jumps:

1. A *pure* context jump is one in which the actual configuration is left unchanged, $Q_{t^+} = Q_{t^-}$, but is thereafter guided by the updated state Ψ_{t^+}.
2. A *mixed* context jump is tied to particle creation or annihilation and can be treated as a combined jump of (Λ_t, Q_t) by using B6 in conjunction with B4.

[26] Here $\widetilde{\mathcal{H}}_{\Lambda_f} := \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$ and $\widetilde{\pi}_{\Lambda_f}(A) := \pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}}$ for $A \in \mathcal{A}_0$ (i.e. an amplification of the representation restricted to the overlap algebra only), for some auxiliary Hilbert space $\mathcal{K}$. $\mathcal{K}$ is purely technical and may be fixed once and for all; see Remark 7.5 below.

7.2 Remarks and Assumptions

Remark 7.1 *(Relation to CBQFT-1) Compared to CBQFT-1, CBQFT-2 makes four structural changes.*

1. *First, the wavefunction no longer lives in a single fixed Fock space: the context Λ_t selects a representation π_{Λ_t} of the field algebra on a Hilbert space $\mathcal{H}_{\Lambda_t}$, and the state vector Ψ_t is understood as an element of $\mathcal{H}_{\Lambda_t}$.*[27] *Whether each of these context-indexed Hilbert spaces are Fock-like*[28] *depends on our choice of Q_t (particles/fields); see Assumption 7.3.*
2. *Second, the Hamiltonians H_{Λ_t} are self-adjoint, so the non-Hermitian "leak" terms used in CBQFT-1 to model absorption are replaced by an explicit Bell-type jump structure for the configuration B4, together with the stochastic context dynamics of B5.*
3. *Third, the context term $H_C(\Lambda)$ of the Hamiltonian, which was invoked in both A2 and A4, has been absorbed into the family $\{H_\Lambda\}_{\Lambda\in\mathcal{C}}$, each acting on its own Hilbert space $\mathcal{H}_\Lambda$.*
4. *Fourth, the context Λ_t itself is promoted from an external control parameter to a stochastic process on $\mathcal{C}$, with context jumps implemented by B5 and B6.*

As we will see presently, these changes restore global norm conservation and full Bohmian equivariance while allowing for feedback between microscopic configurations and macroscopic contexts.

Remark 7.2 *(Local context fields vs. global sectors) At each time t the theory contains a single global wave function $\Psi_t \in \mathcal{H}_{\Lambda_t}$ guiding a single configuration Q_t. The context variable Λ_t may encode a spatially varying field $\Lambda(x,t)$ (for example, different order parameters or detector settings in different regions), so distinct spatial regions can occupy different macroscopic phases at the same time. Such inhomogeneities are nevertheless represented within a single Hilbert space $\mathcal{H}_{\Lambda_t}$ and Hamiltonian H_{Λ_t} at fixed t.*

Remark 7.3 *(Status of Λ_t as a beable) In CBQFT-2 the macroscopic context Λ_t is not introduced as a mere bookkeeping device for unmodelled microscopic degrees of freedom. Rather, Λ_t is stipulated as an additional beable—a physical, coarse-grained macroscopic variable that takes values in $\mathcal{C}$ and enters the fundamental state of the theory alongside the configuration Q_t and the guiding state Ψ_t. Its role is to encode which macroscopic regime is actually realised and to do so in a way that is dynamically efficacious: changes in Λ_t select the active representation π_{Λ_t} and Hamiltonian H_{Λ_t} and thereby redirect the subsequent Bohmian evolution and jump structure.*

Remark 7.4 *(Sector jumps and top–down causation) A change of Hilbert space—and thus of the context Λ_t—is reserved for temporal transitions between unitarily inequivalent sectors (e.g. different thermodynamic phases). Such a transition at time*

[27] Up to the inessential amplification used only to implement context jumps on $\mathcal{A}_0$.

[28] By 'Fock-like' we mean only that $\mathcal{H}_\Lambda$ admits an n-sector grading adequate for Bell-type jump dynamics. Disjointness between π_{Λ_i} and π_{Λ_j} remains allowed.

t_* is modelled as a jump $\Lambda_{t_*^-} \to \Lambda_{t_*^+}$, implemented by the isometry $U_{\Lambda_{t_*^-} \to \Lambda_{t_*^+}}$ of B6. These sector jumps are genuine dynamical events in CBQFT-2: they change H_{Λ_t} (and hence the associated jump rates) and thereby redirect the subsequent Bohmian evolution of Q_t. So macroscopic context changes exert real top–down causal influence. The particle jumps of B4, on the other hand, do not induce a change of representation.

Remark 7.5 *(On the role of the amplification) The auxiliary factor $\mathcal{K}$ in $\widetilde{\mathcal{H}}_{\Lambda_f} := \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$ is introduced only to guarantee the existence of a vector implementation of the same restricted normal state on the overlap algebra $\mathcal{A}_0$ across a context change (see Lemma 7.1). It plays no dynamical role and carries no additional beables. In particular, one does not iterate amplifications along a trajectory: equivalently, one may fix a single sufficiently large $\mathcal{K}$ once and for all and regard each $\widetilde{\mathcal{H}}_\Lambda$ as a notational stabilisation, while all intra-context dynamics and guidance remain entirely within $\mathcal{H}_\Lambda$.*

Assumption 7.1 (Sector disjointness and local overlap) For distinct context values $\Lambda_i \neq \Lambda_j$, the representations $(\pi_{\Lambda_i}, \mathcal{H}_{\Lambda_i})$ and $(\pi_{\Lambda_j}, \mathcal{H}_{\Lambda_j})$ of the field algebra $\mathcal{A}$ are disjoint as *global* representations (i.e. not unitarily equivalent on $\mathcal{A}$), but they are locally quasiequivalent on a chosen overlap algebra $\mathcal{A}_0 \subset \mathcal{A}$ supported in a bounded region away from the macroscopic transition region. In particular, π_{Λ_i} and π_{Λ_j} induce the same normal state space on $\mathcal{A}_0$.

Assumption 7.2 (Configuration in the overlap) For any pair of contexts Λ_i, Λ_j that can be connected by a context jump $\Lambda_i \to \Lambda_j$, we assume there exists a common subalgebra $\mathcal{A}_0 \subset \mathcal{A}$ such that the configuration observable (particle positions or relevant local field configuration) and the associated local density and current operators (entering Axiom B3) belong to $\mathcal{A}_0$.

Assumption 7.3 (Particle contexts) If Q_t in (B1) refers to particle positions, we distinguish a subset $\mathcal{C}_{\mathrm{Fock}} \subseteq \mathcal{C}$ of particle contexts in which the representation $(\pi_\lambda, \mathcal{H}_\lambda)$ admits a Fock-like decomposition into particle-number sectors and supports the particle-configuration observable used in Q_t. Throughout Section 7, we restrict attention to dynamics whose realised contexts satisfy $\Lambda_t \in \mathcal{C}_{\mathrm{Fock}}$, so that the particle guidance law and Bell-type creation/annihilation jumps are well-defined. Contexts outside $\mathcal{C}_{\mathrm{Fock}}$ may be accommodated by switching from particles to fields and adjusting the associated guidance/jump laws (B3 and B4) accordingly; see Section 7.8.

Assumption 7.4 (Well-posedness) For each context value Λ_t and state Ψ_t, the total jump intensity

$$\tau_{\Lambda_t}(Q_t) := \sum_{q' \neq Q_t} \sigma_{\Psi_t, \Lambda_t}(Q_t \to q')$$

is locally integrable, and $t \mapsto H_{\Lambda_t}$ is piecewise strongly continuous. Under these conditions, the piecewise–deterministic Markov process (Ψ_t, Q_t, Λ_t) is well defined.

7.3 Two Consistency Checks

One of the key ways in which CBQFT-2 departs from CBQFT-1 is by allowing the Hilbert space in which the global wave-function is defined to change when a context switch occurs (B5) and the wave function to be re-expressed (B6). We now verify that the axioms of CBQFT-2 fit together coherently: specifically, that the state can be re-expressed consistently across context jumps, and the continuity equation underlying Bohmian guidance (B3) extends globally in time even when the context changes.

7.3.1 Re-Expressing the State Across Context Jumps

When the macroscopic context jumps, $\Lambda_i \to \Lambda_f$, CBQFT-2 re-expresses the state so that all expectations on a common local algebra $\mathcal{A}_0$ are unchanged—that is, the state is transported into the new representation without altering its *local* physical content—while genuinely macroscopic contextual observables (outside $\mathcal{A}_0$) may change. In general, a normal state on $\mathcal{A}_0$ need not admit a vector implementation in the given representation space $\mathcal{H}_{\Lambda_f}$ when $\pi_\Lambda(\mathcal{A}_0)''$ is a type II or type III factor. However, it *can* be implemented after an inessential amplification of the representation that is built into Axiom B6 and discussed in more detail below: $\widetilde{\mathcal{H}}_{\Lambda_f} = \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$. After this, all overlap-algebra observables $A \in \mathcal{A}_0$ are understood to act as $\pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}}$ on $\widetilde{\mathcal{H}}_{\Lambda_f}$. The following lemma indicates how such an isometric re-expression of the state arises from standard assumptions of algebraic QFT.

Lemma 7.1 *(Isometric re-expression from local quasiequivalence) Let Λ_i, Λ_f be two contexts with representations $\pi_{\Lambda_i} : \mathcal{A} \to \mathcal{B}(\mathcal{H}_{\Lambda_i})$ and $\pi_{\Lambda_f} : \mathcal{A} \to \mathcal{B}(\mathcal{H}_{\Lambda_f})$. Fix a C^*-subalgebra $\mathcal{A}_0 \subset \mathcal{A}$, and write $\mathcal{M}_\Lambda := \pi_\Lambda(\mathcal{A}_0)''$ for the generated von Neumann algebra on $\mathcal{H}_\Lambda$. Assume:*

1. *the restrictions of π_{Λ_i} and π_{Λ_f} to $\mathcal{A}_0$ extend to normal representations of $\mathcal{M}_{\Lambda_i}$ and $\mathcal{M}_{\Lambda_f}$;*
2. *π_{Λ_i} and π_{Λ_f} are locally quasiequivalent on $\mathcal{A}_0$, i.e. they induce the same folium of normal states on $\mathcal{A}_0$.*

Fix $\Psi_i \in \mathcal{H}_{\Lambda_i}$ and define a state on $\mathcal{A}_0$ by

$$\omega(A) := \langle \Psi_i, \pi_{\Lambda_i}(A)\Psi_i \rangle, \qquad A \in \mathcal{A}_0.$$

Then there exist an auxiliary Hilbert space $\mathcal{K}$ and a partial isometry

$$U_{\Lambda_i \to \Lambda_f} : \mathcal{H}_{\Lambda_i} \longrightarrow \widetilde{\mathcal{H}}_{\Lambda_f} := \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$$

with initial subspace $\overline{\pi_{\Lambda_i}(\mathcal{A}_0)\Psi_i} \subset \mathcal{H}_{\Lambda_i}$ such that, for all $A \in \mathcal{A}_0$ and all Ψ in that initial subspace,

$$\langle \Psi, \pi_{\Lambda_i}(A)\Psi \rangle = \left\langle U_{\Lambda_i \to \Lambda_f}\Psi,\, \widetilde{\pi}_{\Lambda_f}(A)\, U_{\Lambda_i \to \Lambda_f}\Psi \right\rangle, \tag{25}$$

where $\widetilde{\pi}_{\Lambda_f}(A) := \pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}}$. *In particular,* $U_{\Lambda_i \to \Lambda_f}$ *preserves the norm on its initial subspace, and if* $\Psi_{t^+} := U_{\Lambda_i \to \Lambda_f} \Psi_{t^-}$ *then the restricted state on* $\mathcal{A}_0$ *is preserved across context jumps.*

Proof idea Let $(\mathcal{H}_\omega, \pi_\omega, \Omega_\omega)$ be the GNS triple of the state ω on $\mathcal{A}_0$.

Step 1. Since $\omega(A) = \langle \Psi_i, \pi_{\Lambda_i}(A)\Psi_i \rangle$ on $\mathcal{A}_0$, the GNS representation $(\pi_\omega, \mathcal{H}_\omega)$ is unitarily equivalent to the cyclic subrepresentation of $\pi_{\Lambda_i}|_{\mathcal{A}_0}$ generated by Ψ_i. Hence there is an isometry $V_i : \mathcal{H}_\omega \to \mathcal{H}_{\Lambda_i}$ such that

$$V_i \Omega_\omega = \Psi_i, \qquad V_i \pi_\omega(A) = \pi_{\Lambda_i}(A) V_i \quad \forall\, A \in \mathcal{A}_0.$$

Step 2. By local quasiequivalence (equality of folia), the cyclic representation π_ω is *quasi-contained* in $\pi_{\Lambda_f}|_{\mathcal{A}_0}$ (since equality of folia implies mutual quasi-containment). Equivalently, there exist an auxiliary Hilbert space $\mathcal{K}$ and an isometry $V_f : \mathcal{H}_\omega \to \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$ such that

$$V_f\, \pi_\omega(A) = (\pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}})\, V_f \qquad \forall\, A \in \mathcal{A}_0.$$

The amplification's role is to supply enough multiplicity for such an intertwining isometry to exist, without requiring $\pi_{\Lambda_f}(\mathcal{A}_0)''$ to be a type I factor.

Step 3. Set $U_{\Lambda_i \to \Lambda_f} := V_f V_i^*$. Then $U_{\Lambda_i \to \Lambda_f}$ is a partial isometry with initial subspace $\overline{\pi_{\Lambda_i}(\mathcal{A}_0)\Psi_i}$, and for all $A \in \mathcal{A}_0$ and all Ψ in that subspace,

$$\langle \Psi, \pi_{\Lambda_i}(A)\Psi \rangle = \big\langle U_{\Lambda_i \to \Lambda_f}\Psi,\, (\pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}})\, U_{\Lambda_i \to \Lambda_f}\Psi \big\rangle,$$

which is (25). Norm preservation on the initial subspace follows because V_i and V_f are isometries. □

Remark 7.6 *(On the non-uniqueness of the re-expression map) The isometry* $U_{\Lambda_i \to \Lambda_f} : \mathcal{H}_{\Lambda_i} \to \widetilde{\mathcal{H}}_{\Lambda_f}$ *of Axiom B6 is constrained only by the requirement that it preserve expectation values of observables in the overlap algebra* $\mathcal{A}_0$*. This fixes the post-jump restricted state on* $\mathcal{A}_0$ *(and hence the local density/current data relevant to Bohmian guidance), but it does not determine a unique global map* $U_{\Lambda_i \to \Lambda_f}$*. There are two sources of residual freedom. First, the choice of amplification space* $\mathcal{K}$ *for* $\widetilde{\pi}_{\Lambda_f}(A) = \pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}}$*: different* $\mathcal{K}$ *and different vector implementations of the same normal state on* $\mathcal{A}_0$ *yield identical local expectations. Second, even for fixed* $\mathcal{K}$*, if W is a partial isometry on* $\widetilde{\mathcal{H}}_{\Lambda_f}$ *that commutes with* $\widetilde{\pi}_{\Lambda_f}(\mathcal{A}_0)$ *on the range of* $U_{\Lambda_i \to \Lambda_f}$*, then* $W U_{\Lambda_i \to \Lambda_f}$ *yields the same overlap-algebra statistics. This residual freedom may be regarded as a "gauge": it is invisible on* $\mathcal{A}_0$ *and therefore does not affect the Bohmian guidance field* $v = j/\rho$ *or the resulting trajectories of* Q_t*.*

7.3.2 Global Continuity and Guidance

Although CBQFT-2 permits changes of context—and hence of the underlying Hilbert space—the continuity equation remains valid in a global sense. Recall from Section 2.1 that in ordinary Bell-type QFT the guidance law arises by interpreting the quantum probability current as the flow of an actual configuration distributed according to $|\Psi|^2$. The only new issue in CBQFT-2 is to check that context jumps do not spoil this continuity structure.

Lemma 7.2 *(Global continuity under context jumps) Let* $\{t_k\}_{k\in\mathbb{Z}}$ *be the (at most countable) set of context–jump times, and suppose that on each open interval* (t_k, t_{k+1}) *the context* Λ_t *(and hence the pair* $(\mathcal{H}_{\Lambda_t}, H_{\Lambda_t})$*) is constant.*[29] *Let us assume the following:*

1. *On each open interval* (t_k, t_{k+1})*, Axioms B2 and B3 hold, so that the local density and current defined by* $\rho(x,t)$ *and j(x, t) satisfy the usual continuity equation*

$$\partial_t\rho(x,t) + \nabla\cdot j(x,t) = 0$$

in the classical sense.

2. *At each context jump* t_k*, the isometry of Axiom B6 preserves the expectations of the local density and current operators entering the guidance law,*[30] *so that*

$$\rho(x,t_k^-) = \rho(x,t_k^+), \qquad j(x,t_k^-) = j(x,t_k^+)$$

for all relevant x.[31]

Then the continuity equation

$$\partial_t\rho(x,t) + \nabla\cdot j(x,t) = 0$$

holds globally on spacetime in the sense of distributions. In particular, $v = j/\rho$ *defines a globally consistent Bohmian guidance field for* Q_t *(for* $\rho > 0$*).*

Proof sketch Fix $\phi \in C_c^\infty(\mathbb{R}^{1+3})$ and choose $t_0 < t_N$ with $\mathrm{supp}(\phi) \subset (t_0,t_N)\times\mathbb{R}^3$. Let $t_1 < \cdots < t_{N-1}$ be the finitely many context–jump times in (t_0,t_N). Consider the distributional pairing

$$I := \int_{t_0}^{t_N} dt \int_{\mathbb{R}^3} d^3x\ (\partial_t\rho + \nabla\cdot j)(x,t)\,\phi(x,t).$$

[29] If an amplification is present, H_{Λ_t} is understood to act as $H_{\Lambda_t}\otimes I_{\mathcal{K}}$.

[30] According to Assumption 7.2, the configuration observable and the associated density/current operators lie in the overlap algebra $\mathcal{A}_0$ for any admissible context jump, so their expectations are preserved by the isometry of Axiom B6.

[31] With ρ, j at t_k^+ evaluated using the amplified action $\pi_{\Lambda_f}(\cdot)\otimes I_{\mathcal{K}}$ on $\mathcal{A}_0$ (Lemma 7.1).

To show $I = 0$, we first insert small cutoffs around the jump times: for $\varepsilon > 0$, split the time integral into the union of (i) the bulk intervals $(t_k + \varepsilon, t_{k+1} - \varepsilon)$ and (ii) the ε–neighbourhoods $(t_k - \varepsilon, t_k + \varepsilon)$ of the jump times. On each bulk interval the context is constant, and by assumption (1) $\partial_t \rho + \nabla \cdot j = 0$ classically, so the bulk contribution to I is identically zero for every ε.

Thus I is supported entirely at the jump times. To compute the jump contribution, integrate by parts in t and x on $(t_k - \varepsilon, t_k)$ and $(t_k, t_k + \varepsilon)$. The spatial boundary terms vanish since ϕ has compact spatial support, and the only surviving time–boundary term as $\varepsilon \to 0$ is

$$\int_{\mathbb{R}^3} d^3x \, (\rho(x, t_k^-) - \rho(x, t_k^+)) \, \phi(x, t_k).$$

By assumption (2) the density is continuous across each context jump, i.e. $\rho(\cdot, t_k^-) = \rho(\cdot, t_k^+)$, hence every contribution to I for each k vanishes and we confirm that $I = 0$ as required. Since ϕ was arbitrary, $\partial_t \rho + \nabla \cdot j = 0$ holds globally in the sense of distributions, and wherever $\rho > 0$ the guidance field $v = j/\rho$ is globally consistent across context jumps. □

Remark 7.7 *(Continuity assumptions) Lemma 7.2 only requires continuity of the density and current at context jump times. No assumption is made about the continuity of their time derivatives; the continuity equation is understood in the distributional sense, and the proof uses only the piecewise classical continuity on* (t_k, t_{k+1}) *plus these matching conditions at the jumps.*

7.4 Bohmian Properties

In CBQFT-2 the state triple (Ψ_t, Q_t, Λ_t) evolves according to Axioms B1–B6. Here, we underscore the principal Bohmian dynamical consequences of this structure. First, the guiding state remains normalised *pathwise*. Second, the Bohmian configuration Q_t satisfies full equivariance with respect to $|\Psi_t|^2$.

7.4.1 Norm Conservation

In CBQFT-1, non-Hermitian parts of $H_C(\Lambda)$ lead to norm loss and norm conservation is recovered only at the ensemble level via GKLS dynamics. CBQFT-2 removes these effective leak terms, preventing norm loss.

Proposition 7.1 *(Pathwise norm conservation in CBQFT-2) Under B2 and B6, along almost every sample path of the piecewise–deterministic process* (Ψ_t, Q_t, Λ_t) *one has*

$$\|\Psi_t\| = \|\Psi_{t_0}\| \quad for all t \geq t_0.$$

Proof Fix a realisation of the context process $t \mapsto \Lambda_t$. On any open interval on which Λ_t is constant, Axiom B2 gives unitary dynamics in the active Hilbert space $\mathcal{H}_{\Lambda_t}$, hence $\|\Psi_t\|$ is constant on that interval. At a context jump time t_* with $\Lambda_{t_*^-} = \Lambda_i$ and $\Lambda_{t_*^+} = \Lambda_f$, Axiom B6 supplies a (partial) isometry $U_{\Lambda_i \to \Lambda_f} : \mathcal{H}_{\Lambda_i} \to \widetilde{\mathcal{H}}_{\Lambda_f}$, and we set $\Psi_{t_*^+} = U_{\Lambda_i \to \Lambda_f} \Psi_{t_*^-}$. By definition of a (partial) isometry, $\|\Psi_{t_*^+}\| = \|U_{\Lambda_i \to \Lambda_f} \Psi_{t_*^-}\| = \|\Psi_{t_*^-}\|$. Particle jumps, as specified in Axiom B4, change only the configuration Q_t; Ψ_t is continuous across such jumps and thus its norm is unaffected. Since the full evolution of Ψ_t is obtained by concatenating unitary segments and isometric re-expressions, $\|\Psi_t\|$ remains equal to its initial value $\|\Psi_{t_0}\|$ along each realised trajectory. □

Remark 7.8 *(Local constraints vs. global isometry) The isometry $U_{\Lambda_i \to \Lambda_f}$ is fixed (up to irrelevant gauge freedom) by requiring it to preserve expectation values of local observables in the overlap algebra (in particular, the density and current operators entering the guidance law). This locality of the constraints does not affect global norm conservation: by definition $U_{\Lambda_i \to \Lambda_f}$ is an isometry on its initial subspace, so $\| \Psi_{t_k^+} \| = \| \Psi_{t_k^-} \|$ at each context jump.*

Remark 7.9 *(Improvement over CBQFT-1) In CBQFT-1, in the open-system regime, non-Hermitian context terms model absorption by allowing the single wavefunction Ψ_t to lose norm between jumps, with norm conservation recovered only at the level of the ensemble density matrix ρ_t rather than Ψ_t (Proposition 4.1). CBQFT-2 replaces such effective leak terms by context-dependent Hilbert-space sectors and isometric re-expression (Axiom B6), and particle jumps leave Ψ_t unchanged. The guiding state therefore remains normalised at all times along each path. Norm conservation is thus strictly stronger in CBQFT-2: it holds pathwise, not only after ensemble averaging.*

7.4.2 Equivariance of the Configuration

Equivariance in CBQFT-2 rests on two simple observations. Between context jumps, the pair (Ψ_t, Q_t) evolves as an ordinary Bell-type Bohmian QFT with a self-adjoint Hamiltonian, so the standard $|\Psi_t|^2$-equivariance holds.[32] At a context jump, the configuration Q_t is continuous while the wave function is re-expressed by an isometry that preserves the local density/current data that is relevant to Bohmian guidance. Putting these together yields equivariance for the full CBQFT-2 process.

Proposition 7.2 *(Equivariance of Q_t in CBQFT-2) Assume Axioms B2, B4, and B6. Let $(\mathcal{G}_t)_{t \geq t_0}$ be the (right-continuous) filtration generated by the context process $(\Lambda_s)_{s \leq t}$, i.e.*

$$\mathcal{G}_t := \sigma(\Lambda_s : s \leq t),$$

representing the information carried by the realised context history up to time t. If

[32] Both here and throughout this proof, $|\Psi_t(q)|^2\, dq$ denotes the configuration density in the active particle-context representation, as in Assumption 7.3.

$$\mathbb{P}(Q_{t_0} \in dq \,|\, \mathcal{G}_{t_0}) = |\,\Psi_{t_0}(q)|^2 \, dq \quad a.s.,$$

then for all $t \geq t_0$*,*

$$\mathbb{P}(Q_t \in dq \,|\, \mathcal{G}_t) = |\,\Psi_t(q)|^2 \, dq \quad a.s.$$

In particular, defining the sector-conditional configuration densities

$$\rho_Q^{(i)}(q,t)\, dq := \mathbb{P}(Q_t \in dq \mid \Lambda_t = \Lambda_i), \qquad p_i(t) := \mathbb{P}(\Lambda_t = \Lambda_i),$$

the unconditional density satisfies the mixture identity

$$\rho_Q(q,t) = \sum_i p_i(t)\, \rho_Q^{(i)}(q,t),$$

and by conditional equivariance one has $\rho_Q^{(i)}(q,t) = |\,\Psi_t^{(i)}(q)|^2$ *(a.s.), where* $\Psi_t^{(i)}$ *denotes the random state* Ψ_t *on the event* $\{\Lambda_t = \Lambda_i\}$*.*

Proof Fix $t \geq t_0$ and condition on the realised context history up to time *t*, i.e. on $\mathcal{G}_t$. This determines a partition of $[t_0, t]$ into constant-context intervals $I_k = (t_k, t_{k+1})$ on which $\Lambda_s = \lambda_k$ (with the t_k the context-jump times in $[t_0, t]$).

Step 1. On each I_k the dynamics of (Ψ_s, Q_s) is exactly that of a Bell-type Bohmian QFT with self-adjoint Hamiltonian H_{λ_k} and the corresponding Bell-type jump rates (Axioms B2 and B4). Hence the standard equivariance theorem applies on I_k: if at some time $s_* \in I_k$ one has

$$\mathbb{P}(Q_{s_*} \in dq \,|\, \mathcal{G}_{s_*}) = |\Psi_{s_*}(q)|^2 \, dq,$$

then for all $s \in I_k$,

$$\mathbb{P}(Q_s \in dq \,|\, \mathcal{G}_s) = |\Psi_s(q)|^2 \, dq.$$

Step 2. Let t_k be a context-jump time. By definition, a pure context jump changes Λ but not the configuration, so $Q_{t_k^-} = Q_{t_k^+}$. Meanwhile the state is updated by the isometry,

$$\Psi_{t_k^+} = U_{\Lambda_i \to \Lambda_f} \Psi_{t_k^-}.$$

By Axiom B6, expectations of the local density/current operators entering the guidance law (which lie in the overlap algebra $\mathcal{A}_0$) are preserved across the jump, with post-jump expectations evaluated using the amplified action $\widetilde{\pi}_{\Lambda_f}(A) = \pi_{\Lambda_f}(A) \otimes I_{\mathcal{K}}$ on $\mathcal{A}_0$. In particular, the configuration density $\rho(\cdot, t)$ relevant to guidance/jump rates is unchanged across the jump (a.e. on configuration space). Since $\mathcal{G}_{t_k}$ is right-continuous, it contains the post-jump context value, and therefore the conditional law matches across the jump:

$$\mathbb{P}\left(Q_{t_k^-} \in dq \,|\, \mathcal{G}_{t_k^-}\right) = |\Psi_{t_k^-}(q)|^2\, dq \implies \mathbb{P}\left(Q_{t_k^+} \in dq \,|\, \mathcal{G}_{t_k}\right) = |\Psi_{t_k^+}(q)|^2\, dq.$$

Step 3. Starting from the hypothesis at t_0, apply Step 1 on the first interval and Step 2 at the first jump time, and iterate through all intervals/jumps up to *t*. This yields

$$\mathbb{P}(Q_t \in dq \,|\, \mathcal{G}_t) = |\Psi_t(q)|^2\, dq \quad \text{a.s.}$$

for all $t \geq t_0$. □

Remark 7.10 *(Improvement over CBQFT-1) On each constant-context interval, CBQFT-2 has a self-adjoint generator, so there is no anti-Hermitian part and hence no sink term J_Γ of the kind that appears in Proposition 4.2 when we adopt the open-system regime of CBQFT-1. Equivariance is therefore restored in the pathwise sense with respect to $|\Psi_t|^2$, rather than only at the ensemble level via a GKLS state in an effective open-system regime.*

7.5 Contextual Properties

CBQFT-2 assigns two roles to the context variable Λ_t. First, it labels the active global representation $(\pi_{\Lambda_t}, \mathcal{H}_{\Lambda_t})$, and the disjointness of representations prevents normal coherent superpositions *between* distinct macroscopic contexts. Second, Λ_t evolves by a stochastic jump law driven by microscopic quantities. Together, these features yield (i) exact decoherence of the context sectors at the algebraic level, and (ii) an intrinsically time-directed, coarse-grained Markov dynamics for context changes. Let us discuss these features.

7.5.1 Decoherence and Effective Classicality of Context

In CBQFT-1, the context variable Λ is treated as an externally fixed classical parameter; decoherence of pointer states is imported from standard open-system arguments applied to Ψ_t and its environment [7, 21, 46, 48, 60]. In CBQFT-2, by contrast, the context Λ_t is itself a stochastic dynamical variable, and Assumption 7.1 places macroscopically distinct contexts in disjoint *global* representations of the quasilocal algebra $\mathcal{A}$, while allowing local quasiequivalence on overlap algebras $\mathcal{A}_0$ supported away from the transition region. This yields an *exact* decoherence (that is, an enforced superselection) between macroscopically distinct contexts at the algebraic level.

Proposition 7.3 *(Exact decoherence between context sectors) Under Assumption 7.1 and Axioms B1–B5, the induced ensemble (algebraic) state ω_t on $\mathcal{A}$ is classical in the context label. More precisely, ω_t is a convex mixture of normal sector states $\omega_t^{(i)}$ supported in the representations $(\pi_{\Lambda_i}, \mathcal{H}_{\Lambda_i})$:*

$$\omega_t = \sum_i p_i(t)\, \omega_t^{(i)}, \qquad p_i(t) := \mathbb{P}(\Lambda_t = \Lambda_i).$$

Proof idea Fix t. For each context value Λ_i, let $p_i(t) := \mathbb{P}(\Lambda_t = \Lambda_i)$. Define the conditional microscopic state in sector $\mathcal{H}_{\Lambda_i}$ by

$$\rho_t^{(i)} := \mathbb{E}\big(\,|\Psi_t\rangle\langle\Psi_t|\,\big|\,\Lambda_t = \Lambda_i\big)\,.$$

Since $\rho_t^{(i)}$ is a positive trace-class operator on $\mathcal{H}_{\Lambda_i}$ (or on the trivially amplified space $\mathcal{H}_{\Lambda_i} \otimes \mathcal{K}$), it induces a normal algebraic state $\omega_t^{(i)}$ on $\mathcal{A}$ via

$$\omega_t^{(i)}(A) := \mathrm{tr}\big(\rho_t^{(i)}\,\pi_{\Lambda_i}(A)\big), \qquad A \in \mathcal{A}.$$

By conditioning on the discrete classical variable Λ_t and using the law of total expectation, the unconditional ensemble state is therefore

$$\omega_t(A) = \sum_i p_i(t)\,\omega_t^{(i)}(A), \qquad A \in \mathcal{A},$$

i.e. $\omega_t = \sum_i p_i(t)\,\omega_t^{(i)}$. Because the representations $(\pi_{\Lambda_i}, \mathcal{H}_{\Lambda_i})$ and $(\pi_{\Lambda_j}, \mathcal{H}_{\Lambda_j})$ are disjoint for $i \neq j$ (Assumption 7.1), there is no normal state of $\mathcal{A}$ whose GNS representation contains coherent superpositions across distinct context Hilbert spaces.[33] Equivalently, normal states of $\mathcal{A}$ admit only sector-wise (block-diagonal) decompositions, with no normal interference terms between distinct $\Lambda_i \neq \Lambda_j$. This mirrors the standard picture of superselection sectors and disjoint representations in algebraic QFT [18, 25, 45]. □

Remark 7.11 *(Beable-level mixture) Since Λ_t is a discrete classical label in this formalisation of CBQFT-2, the joint law of (Λ_t, Q_t) always disintegrates with respect to Λ_t. Writing $\rho_Q^{(i)}(q,t)\,dq := \mathbb{P}(Q_t \in dq \mid \Lambda_t = \Lambda_i)$, one has*

$$\mathbb{P}(\Lambda_t \in d\lambda, Q_t \in dq) = \sum_i \delta_{\Lambda_i}(d\lambda)\,p_i(t)\,\rho_Q^{(i)}(q,t)\,dq.$$

This identity is not an additional dynamical assumption; it is included only to make explicit that the primitive ontology statistics are also a classical mixture in the context label, mirroring the block-diagonal sector decomposition of ω_t.

Remark 7.12 *(Improvement over CBQFT-1) CBQFT-1 stipulates a classical context variable Λ but does not explain, within the theory, why macroscopically distinct contexts behave as mutually non-interfering "pointer alternatives". CBQFT-2 improves on this in two ways. First, the context Λ_t is dynamical: its jump Markov evolution is explicitly linked to microscopic quantities via Axiom B5. Second, Assumption 7.1*

[33] When an amplification $\widetilde{\mathcal{H}}_\Lambda = \mathcal{H}_\Lambda \otimes \mathcal{K}$ is used, we understand it as the trivial tensor extension of the full representation, $\widetilde{\pi}_\Lambda(A) := \pi_\Lambda(A) \otimes I_{\mathcal{K}}$ for all $A \in \mathcal{A}$, even though only overlap-algebra expectations on $\mathcal{A}_0$ are constrained by Axiom B6.

places macroscopically distinct contexts in disjoint global sectors, while allowing local quasiequivalence on overlap algebras $\mathcal{A}_0$. As a result, the global state is block-diagonal in the context label: there is no normal interference between distinct macroscopic contexts, and $t \mapsto \Lambda_t$ functions as an internally generated classical record of the macroscopic state of the apparatus and environment.

Remark 7.13 *(Microscopic entanglement vs. macroscopic classicality) Proposition 7.3 constrains only coherences between distinct context sectors. Once the global state is block-diagonal in Λ_t, there is no normal quantum coherence linking different values of Λ_t, but within each sector $(\pi_{\Lambda_i}, \mathcal{H}_{\Lambda_i})$ the state Ψ_t may still exhibit arbitrary microscopic entanglement between local observables. This entanglement is propagated by the unitary dynamics generated by H_{Λ_t} and transported by the isometries $U_{\Lambda_i \to \Lambda_f}$ across context changes. Thus CBQFT-2 enforces classicality of the context record at the sector level without eliminating fine-grained quantum correlations within each context sector, leaving open the familiar (typically inexact) decoherence mechanisms that may operate within each sector [46, 48, 60].*

7.5.2 Irreversibility of the Context Process

CBQFT-2 provides a natural setting in which to discuss the *irreversibility* of the macroscopic domain. In CBQFT-1 the context variable Λ is externally prescribed, so any time-asymmetry in its behaviour is imported. By contrast, in CBQFT-2 the context Λ_t is a stochastic dynamical variable with transition rates $W_{\Lambda_i \to \Lambda_j}(Q_t, \mathcal{J}_t)$ (Axiom B5), driven by microscopic quantities such as currents or absorbed flux. This yields a notion of entropy production that quantifies an intrinsic arrow of time for the context dynamics.

Proposition 7.4 *(Effective Markov description and arrow of time) Once again, let $p_i(t) := \mathbb{P}(\Lambda_t = \Lambda_i)$. Then $p_i(t)$ satisfies a closed (Kolmogorov forward) master equation for a continuous-time Markov chain on the label set $\{\Lambda_i\}$, with time-dependent effective transition rates $\bar{W}_{i \to j}(t)$ obtained by conditional averaging of the microscopic kernel $W_{\Lambda_i \to \Lambda_j}(Q_t, \mathcal{J}_t)$. Moreover, the effective context dynamics admits a canonical entropy production rate $\dot{S}_{\mathrm{prod}}(t) \geq 0$ (in the sense of stochastic thermodynamics). Whenever detailed balance is violated over an interval, the context process exhibits an intrinsic arrow of time at the coarse-grained level: $\dot{S}_{\mathrm{prod}}(t) > 0$.*

Proof sketch Define $p_i(t) := \mathbb{P}(\Lambda_t = \Lambda_i)$. Starting from the Markov jump dynamics of Λ_t with hazards $W_{\Lambda_i \to \Lambda_j}(Q_t, \mathcal{J}_t)$, one obtains, by conditioning on $\{\Lambda_t = \Lambda_i\}$ and averaging over the microscopic ensemble, the effective rates

$$\bar{W}_{i \to j}(t) := \mathbb{E}\big[\, W_{\Lambda_i \to \Lambda_j}(Q_t, \mathcal{J}_t) \,\big|\, \Lambda_t = \Lambda_i \,\big].$$

Marginalising over (Q_t, Ψ_t) then yields the Kolmogorov forward equation

$$\dot{p}_i(t) \;=\; \sum_j \big[\bar{W}_{j\to i}(t)\, p_j(t) - \bar{W}_{i\to j}(t)\, p_i(t)\big],$$

i.e. a (generally time-inhomogeneous) continuous-time Markov chain on $\{\Lambda_i\}$. For such Markov jump processes, stochastic thermodynamics associates the entropy production rate

$$\dot{S}_{\rm prod}(t) := \frac{1}{2}\sum_{i,j}\Big[\bar{W}_{i\to j}(t)p_i(t) - \bar{W}_{j\to i}(t)p_j(t)\Big]\log\frac{\bar{W}_{i\to j}(t)p_i(t)}{\bar{W}_{j\to i}(t)p_j(t)},$$

which is nonnegative by the elementary inequality $(a-b)\log(a/b) \geq 0$ for $a, b > 0$ (with the usual continuous extension when one of the terms vanishes). Moreover, $\dot{S}_{\rm prod}(t) = 0$ iff $\bar{W}_{i\to j}(t)p_i(t) = \bar{W}_{j\to i}(t)p_j(t)$ for all i, j, which is precisely the (instantaneous) detailed-balance condition with respect to a stationary reference distribution when such a reference exists. Thus, whenever detailed balance is violated, $\dot{S}_{\rm prod}(t) > 0$ and the coarse-grained context dynamics carries a thermodynamic arrow of time.[34] □

Remark 7.14 *(Arrow of time for context) CBQFT-2 thus secures an intrinsic notion of irreversibility for macroscopic context changes. The exact decoherence between sectors (Proposition 7.3) makes $t \mapsto \Lambda_t$ a genuinely classical record of the macroscopic state, while Axiom B5 supplies a Markovian dynamics for that record, with entropy production determined by its coupling to the microscopic Bohmian degrees of freedom. CBQFT-1 can mimic such an arrow only by stipulating a time-directed protocol for $\Lambda(t)$; CBQFT-2 embeds the asymmetry in the joint dynamics of (Ψ_t, Q_t, Λ_t) itself.*

7.6 Limiting Regimes

The following three limits locate CBQFT-2 relative to familiar cases.

- **Fixed context (BQFT-limit):** $W_{i\to j} \equiv 0$ for all $i \neq j$. Then $\Lambda(t) \equiv \Lambda_0$ and the model reduces to a Bell–type Bohmian QFT in the single representation $\mathcal{H}_{\Lambda_0}$ (discussed in Section 2).
- **Exogenous context (CBQFT-1, Hermitian regime):** $W_{i\to j}$ independent of (Q_t, Ψ_t) (no feedback). Context modulates the in-sector dynamics but is not updated by micro-events; this reproduces CBQFT-1 in the Hermitian regime (Section 4).
- **Coarse-grained absorption (CBQFT-1, open-system regime):** In addition to exogenous context, pass to a description in which Λ (and any additional absorption/environment degrees of freedom it stands for) is not tracked explicitly. The resulting unconditional ensemble evolution is GKLS, and the corresponding conditional branch is generated by an effective non-Hermitian Hamiltonian:

[34] For background and standard variants of these constructions, see Schnakenberg [49], Lebowitz–Spohn [35], and Seifert [51].

$H_{\text{eff}} = H - \frac{i}{2}\Gamma$.

7.7 Relation to Hybrid Classical–Quantum Models

A variety of "hybrid" approaches to quantum theory have been proposed that seek to couple a classical system to a quantum one.[35] CBQFT-2 fits into this broad family, but in a particularly simple and well-controlled way. The macroscopic context is treated as a *classical label* $\lambda \in \mathcal{C}$, while the microscopic degrees of freedom form a quantum field system. λ can encode macroscopic geometric and boundary information (e.g. cavity geometry or curvature data), while the quantum sector evolves as a QFT on that classical background.[36] Let us fix once and for all an auxiliary Hilbert space $\mathcal{K}$ (chosen so that it is "large enough" to support the amplifications required by Axiom B6 and Lemma 7.1), and define the stabilised sector spaces

$$\widetilde{\mathcal{H}}_\lambda := \mathcal{H}_\lambda \otimes \mathcal{K}, \qquad \widetilde{\pi}_\lambda(A) := \pi_\lambda(A) \otimes I_{\mathcal{K}} \quad (A \in \mathcal{A}_0).$$

Then one may package the hybrid classical–quantum kinematics as

$$\widehat{\mathcal{H}} = \bigoplus_{\lambda \in \mathcal{C}} \big(\widetilde{\mathcal{H}}_\lambda \otimes \mathbb{C}|\lambda\rangle\big),$$

where $\{|\lambda\rangle\}$ is an orthonormal basis for the macroscopic context register. Since the register is treated as classical, we restrict attention to *classical–quantum* (block-diagonal) states on $\widehat{\mathcal{H}}$ of the form

$$\widehat{\rho} = \sum_{\lambda \in \mathcal{C}} p_\lambda \, \widetilde{\rho}_\lambda \otimes |\lambda\rangle\langle\lambda|, \qquad p_\lambda \geq 0, \ \sum_\lambda p_\lambda = 1,$$

which may be read as: "the context has value λ with probability p_λ, and conditioned on that, the stabilised quantum state is $\widetilde{\rho}_\lambda$ on $\widetilde{\mathcal{H}}_\lambda$."[37] There are correspondingly two kinds of evolution, already identified in Section 7.1.

1. **Within a fixed context** ($\Lambda_t = \lambda$): the stabilised state in $\widetilde{\mathcal{H}}_\lambda$ evolves unitarily under the self-adjoint generator $\widetilde{H}_\lambda := H_\lambda \otimes I_{\mathcal{K}}$ (Axiom B2).
2. **Context switches** ($\lambda \to \mu$): at random jump times the classical label changes according to the kernel of Axiom B5, and the quantum state is re-expressed by an isometry $\widetilde{U}_{\lambda\to\mu} : \widetilde{\mathcal{H}}_\lambda \to \widetilde{\mathcal{H}}_\mu$, chosen so as to preserve overlap expectations $\langle\Psi, \widetilde{\pi}_\lambda(A)\Psi\rangle$ for $A \in \mathcal{A}_0$ (Axiom B6).

[35] See for instance the bracket-based and ensemble approaches of Peres and Terno, Hall and Reginatto [26, 41]. For a more recent stochastic hybrid, see the model motivated by classical gravity due to Oppenheim and collaborators [38].

[36] This is consonant with the sensitivity of vacuum structure to boundaries and geometry, as in Casimir-type settings: [24, 32, 52, 56].

[37] Equivalently, one may take the algebra of observables to be block-diagonal in the $\{|\lambda\rangle\}$ basis, so that coherences between distinct labels are superselected/unobservable.

Viewed on $\widehat{\mathcal{H}}$, each realised context switch $\lambda \to \mu$ induces a completely positive, trace-preserving update on the stabilised quantum block,

$$\widetilde{\rho}_\lambda \longmapsto \widetilde{\mathcal{E}}_{\lambda\to\mu}(\widetilde{\rho}_\lambda) := \widetilde{U}_{\lambda\to\mu}\,\widetilde{\rho}_\lambda\,\widetilde{U}^*_{\lambda\to\mu}, \qquad \mathrm{tr}\,\widetilde{\mathcal{E}}_{\lambda\to\mu}(\widetilde{\rho}_\lambda) = \mathrm{tr}\,\widetilde{\rho}_\lambda.$$

Averaging over the stochastic context outcomes (with probabilities supplied by Axiom B5) yields a standard CPTP evolution on the block-diagonal classical–quantum states.

By contrast, many other hybrid models seek to couple a continuous classical phase space (x, p) to a quantum system via modified commutator brackets or through mean-field terms [41], which make it difficult to maintain consistency. Typical problems include nonconservation of energy, the absence of any genuine back-reaction of the quantum system on the classical one, negative probabilities, and an incorrect classical limit in the absence of interaction (for a succinct overview, see [26, pp. 1–2]). CBQFT-2 sidesteps many of these difficulties; partly because its classical component is treated here as a discrete register rather than a continuum phase space, and partly because the quantum dynamics within each context sector is unmodified and unitary.

The configuration-ensemble approach of Hall and Reginatto, on the other hand, also seeks to avoid these problems, but it does so by treating the joint classical–quantum system as a single ensemble on a joint configuration space, described by a probability density P and a canonically conjugate field S defined on this space. A suitable composite ensemble Hamiltonian $H[P, S]$ then generates coupled evolution equations for P and S that are carefully engineered to ensure conservation of probability and energy, the presence of back-reaction, and the correct classical equations of motion in the limit of no interaction for a wide class of interactions.

CBQFT-2 takes a different route to this hybrid approach too: (i) its classical part (at least, in this formalisation) is a discrete register encoding the macroscopic context, not a continuum of phase-space variables nor an ensemble over a continuous classical configuration space; (ii) the microscopic degrees of freedom always form an ordinary (Bohmian) QFT in each sector, so that no modification of the underlying quantum or Bohmian dynamics is required; (iii) quantum evolution inside each sector is strictly unitary, and randomness in the hybrid register–sector dynamics enters through context jumps (in addition to the Bell-type stochastic particle jumps); and (iv) back-reaction is handled by the context transition rates $W_{\lambda\to\mu}(Q_t, \mathcal{J}_t)$, rather than by continuous mean-field forces or by a single canonical Hamiltonian flow on a joint classical–quantum configuration ensemble. In particular, we do not introduce additional conjugate variables for the classical context or a joint ensemble Hamiltonian $H[P, S]$; instead, the hybrid structure is encoded in the family $\{H_\lambda, U_{\lambda\to\mu}\}$ together with the stochastic dynamics of Λ_t.

This has two pragmatic advantages over the configuration-ensemble approach. First, CBQFT-2 can be implemented directly on any Bohmian QFT model simply by specifying the context-labelled Hamiltonians, isometries, and context kernel, without having to construct and analyse a full classical–quantum ensemble Hamiltonian on the combined configuration space. Second, on the extended space $\widehat{\mathcal{H}}$ the ensemble dynamics is manifestly CPTP by construction, so issues of positivity and norm conservation reduce to the familiar structural conditions on quantum channels, rather than to functional-analytic properties of a continuum ensemble Hamiltonian. From this point of view, CBQFT-2 provides a particularly transparent example of a consistent hybrid classical–quantum dynamics that is tailored to Bohmian QFT.

7.8 Summary

The second version of our model, CBQFT-2, formalises context in Bohmian QFT by embedding the quantum state Ψ_t, particle configuration Q_t, and context Λ_t into a unified dynamical framework. Unlike CBQFT-1, which treats context as an externally prescribed background, CBQFT-2 models Λ_t as a dynamical, piecewise-constant stochastic variable. Each value of Λ_t selects a representation π_{Λ_t} of the field algebra $\mathcal{A}$ on a Hilbert space $\mathcal{H}_{\Lambda_t}$.

Within each context sector, Ψ_t evolves unitarily under a self-adjoint Hamiltonian H_{Λ_t}, preserving the norm and sustaining well-defined Bohmian guidance laws for Q_t. Contextual changes are governed by a probabilistic kernel $P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t)$, which specifies how macroscopic transitions—such as detection events or phase transitions—trigger jumps between representation sectors. Upon such transitions, the state Ψ_t is re-expressed in the new Hilbert space $\mathcal{H}_{\Lambda_f}$ by an isometry, so that the relevant algebraic expectations are preserved.

7.8.1 Benefits

This formalism achieves several goals:

- It secures *pathwise* norm conservation and full Bohmian equivariance across both unitary evolution and stochastic jumps.
- It provides a clean account of how measurement-like updates, decoherence between macroscopic contexts, and an arrow of time for the context process can emerge from context transitions and their Markovian dynamics.
- It realises a feedback loop of micro–macro causation: quantum events influence the macroscopic context Λ_t, and in turn Λ_t governs the dynamics of both Ψ_t and Q_t via H_{Λ_t} and jump rates.
- It generalises and subsumes earlier models such as CBQFT-1 and Bell-type Bohmian QFTs as limiting cases (e.g. when context is fixed or trivial, or when feedback from Q_t to Λ_t is switched off).

In essence, CBQFT-2 offers a Bohmian QFT with dynamically structured Hilbert spaces, allowing physical context to evolve in response to particle-level events. The use of (typically) inequivalent representations to model macroscopic distinctions lends the theory a form of structural realism about context, aligned with insights from algebraic QFT [25, 29, 45]. The contextual dynamism is not imposed ad hoc, but built into the formal machinery of the theory itself, via the coupled evolution of (Ψ_t, Q_t, Λ_t), algebraic states, and stochastic transition rules. In doing so, CBQFT-2 offers a unified framework in which Bohmian dynamics, statistical irreversibility of macroscopic context, and measurement-like phenomena are placed on the same footing.

7.8.2 Limitations

Nonetheless, the scope of this formulation has limitations. For the sake of clarity, and in order to keep close contact with the standard Bell-type Bohmian QFT with which we began in Section 2, we have been working in a *particle ontology* regime and restricting our attention accordingly to a subclass of contexts $\mathcal{C}_{\text{Fock}} \subseteq \mathcal{C}$ for which the associated representations admit a Fock-like particle-number grading, so that a particle configuration Q_t and Bell-type creation/annihilation rates are well-defined on each constant-context interval. Throughout this section we assumed $\Lambda_t \in \mathcal{C}_{\text{Fock}}$.

This restriction can be lifted if one wishes to treat genuinely non-Fock contexts (e.g. thermal/KMS sectors). One may replace the particle beables in Axiom B1 by a field configuration or other local beable defined on the overlap algebra $\mathcal{A}_0$, or adopt a composite primitive ontology of particles *and* fields, and formulate guidance directly in terms of the corresponding local densities and currents in each sector. In such a generalisation, Axiom B4 would be modified or omitted (or restricted to those sectors admitting a particle-number grading), while the context dynamics (Axiom B5) and isometric re-expression principle (Axiom B6) continue to play the same role.

8 CBQFT-2: Implications

Having discussed the axioms of CBQFT-2, we now draw out some of its main conceptual and physical consequences. As we shall see, once macroscopic context is treated as something real and dynamical—selecting the active representation and being updated by microscopic events—a number of familiar phenomena can be analysed as internal features of the theory.

8.1 Macro-to-Micro Information Flow in CBQFT-2

CBQFT-2 reconfigures macro-to-micro information flow by embedding context not as an external parameter but as a dynamically evolving part of the ontology. Whereas in CBQFT-1 the context Λ modulates the dynamics through a fixed Hamiltonian term $H_C(\Lambda)$ on a single Fock space, in CBQFT-2 context is encoded via the choice of a (typically unitarily inequivalent) representation π_{Λ_t} of the quantum field algebra $\mathcal{A}$.

Feedback is now built into the architecture. Particle configurations Q_t and their histories influence the stochastic evolution of Λ_t via the kernel $W_{\Lambda_i \to \Lambda_f}(Q_t, \mathcal{J}_t)$; the value of Λ_t in turn selects the active representation $(\pi_{\Lambda_t}, \mathcal{H}_{\Lambda_t})$ and Hamiltonian H_{Λ_t} that govern the evolution of Ψ_t; and Ψ_t then determines the Bohmian guidance and jump processes for Q_t.[38] This closes the causal loop between the microscopic and the macroscopic without reducing the latter to the former: the system's macroscopic form influences the Bohmian particle trajectories, and those micro-trajectories feed back into the dynamics of form.

[38] Strictly, at a context jump $\Lambda_i \to \Lambda_f$ the re-expression map may land in an inessential amplification $\widetilde{\mathcal{H}}_{\Lambda_f} := \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$, with overlap observables acting as $\pi_{\Lambda_f}(\cdot) \otimes I_{\mathcal{K}}$ on $\mathcal{A}_0$; we suppress this stabilisation factor except when needed.

In this sense CBQFT-2 realises a robust kind of top-down causation. Unlike in CBQFT-1, where context was statically imposed from outside, CBQFT-2 shows how macroscopic order can both emerge from and constrain micro-behaviour. The theory does not rely on extra collapse postulates or heuristic boundary conditions; instead, it internalises context as a dynamically evolving part of the quantum-field ontology. Information flow becomes a dynamical, algebraically grounded feature of the theory itself, providing a unified framework in which measurement, temporal asymmetry, and ontological commitment can be analysed as aspects of a single evolving structure.

8.2 Measurement and the Arrow of Time in CBQFT-2

Having characterised the bidirectional coupling between Λ_t, Ψ_t, and Q_t, we can now turn to a central application: the analysis of measurement and the emergence of an arrow of time. In CBQFT-2, measurement events can be naturally interpreted as *context transitions*: stochastic jumps $\Lambda_i \to \Lambda_f$ governed by the transition kernel $P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t)$ of B5, where $\mathcal{J}_t$ encodes relevant particle-level data such as detector activation, local currents, or interaction histories. Such transitions correspond to shifts in the representational structure of the quantum system: distinct values of Λ_t index (typically) inequivalent Hilbert spaces $\mathcal{H}_{\Lambda_t}$ and representations π_{Λ_t}.

A measurement outcome is thus associated with a re-indexing of the field representation sector—a structural update that reflects a change in the system's macroscopic organisation. In particular, in CBQFT-2 the outcome is not identified merely with a "pointer-shaped" configuration of particles (as it is often presented in standard Bohmian measurement theory), but with the realised value of the macroscopic context register Λ_t (and hence the active sector $(\pi_{\Lambda_t}, \mathcal{H}_{\Lambda_t})$) to which the microscopic configuration Q_t is coupled. The pointer configuration may still serve as a concrete microphysical realisation of that record, but the record itself is encoded at the level of context.

Crucially, this update does not amount to a fundamental, discontinuous "collapse" of the wavefunction, as in standard quantum mechanics or GRW/CSL theory (see Section 9 for comparisons). Instead, at a context jump $\Lambda_i \to \Lambda_f$ (Axiom B5) the state is re-expressed isometrically (Axiom B6) in such a way that all expectations on the overlap algebra $\mathcal{A}_0$ are preserved (and hence, in particular, the local density/current data relevant to guidance). In this sense the *local* physical content (outside the apparatus region) is unchanged at the instant of the jump, while genuinely macroscopic observables associated with the apparatus (e.g. detector currents) can change their expectations because they are not elements of $\mathcal{A}_0$.

The same formalism also underwrites a natural arrow of time. Within each sector $\mathcal{H}_{\Lambda_t}$ the evolution generated by H_{Λ_t} is unitary and time-reversal symmetric; likewise, the isometries $U_{\Lambda_i \to \Lambda_f}$ can in principle be inverted on their ranges. Temporal asymmetry does not arise from the microscopic dynamics of Ψ_t and Q_t taken in isolation, but from the stochastic context process Λ_t and its statistical properties. In physically relevant situations the transition rates $W_{\Lambda_i \to \Lambda_f}(Q_t, \mathcal{J}_t)$ are biased towards macroscopically amplifying, entropy-increasing processes, such as detection, amplification, and thermalisation, whereas reverse transitions would require finely tuned control of many environmental degrees of freedom. As shown in Section 7.5.2, this asymmetry of the context kernel gives rise to a positive entropy production along typical histories of Λ_t, and hence to a thermodynamic arrow of time.

In this way, CBQFT-2 interprets both measurement and time's arrow as consequences of its context-sensitive architecture. Measurement corresponds to a context transition driven by microscopic dynamics and recorded in the macroscopic sector label Λ_t; effective collapse and decoherence emerge from the block structure induced by these transitions. Time-asymmetry arises from the irreversible statistics of the context process, shaped by thermodynamic and environmental constraints, rather than from any fundamental breakdown of unitary (or isometric) quantum evolution. This unifies the phenomenology of collapse, irreversibility, and macro-objectivity within a single, algebraically grounded dynamical framework.

8.3 Causality and Ontology in CBQFT-2

The same feedback architecture has direct implications for how we understand causality and ontology in CBQFT-2. In particular, it clarifies in what sense the macroscopic context is both dynamically active and ontologically structured. Unlike CBQFT-1, CBQFT-2 supports a genuine *two-way* causal coupling between micro- and macro-level quantities. On the one hand, the macroscopic context Λ_t influences the evolution of Ψ_t and Q_t through the context-indexed Schrödinger equation and the corresponding guidance and jump laws (Axioms B2–B4), creating and shaping the Bohmian trajectories. On the other hand, particle-level events—encoded in the configuration Q_t together with the coarse-grained history $\mathcal{J}_t$—enter the context kernel $W_{\Lambda_i \to \Lambda_j}(Q_t, \mathcal{J}_t)$ and can trigger stochastic transitions of Λ_t itself (Axiom B5). A detection event or a symmetry-breaking fluctuation can therefore induce a context switch, after which a *new* in-sector Hamiltonian H_{Λ_t} (and hence a new guidance field and Bell-type jump structure) becomes operative.

This micro–macro coupling is not intended as a mere redescription of the microscopic particle configuration (or the quantum fields). In CBQFT-2 the context Λ_t is stipulated as an additional *primitive* beable: a macroscopic register that forms part of the fundamental state of the theory alongside Q_t and Ψ_t. Its ontological role is to encode which macroscopic regime is actually realised at any given moment and to do so in a dynamically efficacious way: changes in Λ_t select the active representation π_{Λ_t} and Hamiltonian H_{Λ_t} and thereby redirect the subsequent Bohmian motion and the admissible configuration-space transitions.

To avoid ontological inflation, CBQFT-2 introduces a natural equivalence relation on the space of context labels. For $\Lambda_1, \Lambda_2 \in \mathcal{C}$ we write $\Lambda_1 \sim \Lambda_2$ iff the corresponding representations $\pi_{\Lambda_1}, \pi_{\Lambda_2}$ of the quasilocal algebra $\mathcal{A}$ are *unitarily equivalent* (globally), i.e. there exists a unitary $U : \mathcal{H}_{\Lambda_1} \to \mathcal{H}_{\Lambda_2}$ such that $U\,\pi_{\Lambda_1}(A)U^{-1} = \pi_{\Lambda_2}(A)$ for all $A \in \mathcal{A}$. This yields the quotient $\mathcal{C}_{\mathrm{phys}} := \mathcal{C}/\sim$, whose elements $[\Lambda] \in \mathcal{C}_{\mathrm{phys}}$ label the physically distinct context sectors of the theory. The full label space $\mathcal{C}$ may encode fine-grained microstructural detail, modelling convenience, or gauge-dependent choices, but only *globally* inequivalent representation sectors are treated as metaphysically distinct.

This is compatible with Axiom B6. Isometric re-expression at context jumps relies only on *local* quasiequivalence on a chosen overlap algebra $\mathcal{A}_0 \subset \mathcal{A}$, and thus does not require global unitary equivalence of the full representations. Moreover, in non–type I situations the re-expression may require an inessential amplification on the target side: $\widetilde{\mathcal{H}}_{\Lambda_f} = \mathcal{H}_{\Lambda_f} \otimes \mathcal{K}$, with $\mathcal{A}_0$ acting as $\pi_{\Lambda_f}(\cdot) \otimes I_{\mathcal{K}}$ (Lemma 7.1). This does not change the ontology of contexts: $\mathcal{K}$ carries no additional beables and is introduced only to guarantee a vector implementation of the same restricted normal state on $\mathcal{A}_0$. The quotient $\mathcal{C}_{\rm phys}$ therefore captures the intended ontology (macroscopically distinct sectors), while the isometry principle captures the intended physics (preservation of local content across a context switch).

The appeal to (typically) unitarily inequivalent representations in individuating $\mathcal{C}_{\rm phys}$ articulates a structural fact familiar from algebraic QFT: in systems with infinitely many degrees of freedom, genuinely distinct macroscopic phases are naturally modelled by globally disjoint representations of the same quasilocal algebra [25, 29, 45]. In CBQFT-2, however, the direction of dependence is $\Lambda_t \mapsto (\pi_{\Lambda_t}, \mathcal{H}_{\Lambda_t})$: Λ_t is the primitive label of the active sector, and the representation-theoretic machinery is the QFT-level means of *encoding* that macroscopic fact. Nothing here requires positing, in addition to Q_t and Λ_t, an infinite-dimensional configuration beable corresponding to the field modes; rather, the algebraic structure is doing representational work for the dynamics and for the individuation of macroscopic regimes. The intended ontology is thus in the spirit of the primitive ontology (PO) approach to quantum theory [2]: just as PO does not treat the wave function Ψ_t as an additional material beable (even though it lives on a high-dimensional space and carries enormous structure), CBQFT-2 likewise does not reify the "infinite modes" of the quantum fields as beables. They are part of the mathematical machinery that encodes the dynamical role of the non-PO variables.

With these qualifications in place, a hylomorphic reading of CBQFT-2 becomes available (cf. [31]): each $[\Lambda] \in \mathcal{C}_{\rm phys}$ can be viewed as a *form* (or formal constraint) in an Aristotelian sense—an organising principle that fixes stable behavioural dispositions of the micro-ontology through the sector-indexed dynamics [30, 53]. These forms do not merely "modulate" a fixed background of laws. They enter into the concrete realisation of observables (via π_Λ), the available Hamiltonians $\{H_\Lambda\}$, and thereby into which effective particle species, configuration sectors, and creation–annihilation channels are available in a given macroscopic regime. On this construal, form is causally efficacious in a controlled way: it acts through the sector-indexed laws and the admissible transitions between sectors.[39]

Given this integration of Aristotelian forms within the quantum-field ontology, it is natural to introduce an alternative label for CBQFT-2: we can call it *Hylomorphic Bohmian Quantum Field Theory*. This is not intended as a new formalism but as a philosophical gloss that highlights the core innovation: macroscopic context is no longer a passive background parameter but a dynamical, ontologically basic register that both shapes and is shaped by quantum events. The isometric re-expression on overlap algebras then ensures that this bidirectional interplay remains compatible

[39] See [55] for a discussion of this hylomorphic ontology applied to a simplified version of CBQFT, and [54] for an application of hylomorphism to standard Bohmian theory.

with the usual locality constraints for all observables in $\mathcal{A}_0$ (and, in particular, with the no-signalling behaviour of local measurement statistics); our next concern.

8.4 Non-Locality and Relativity in CBQFT-2

We must finally consider how this context-sensitive, hylomorphic structure sits with quantum non-locality and relativistic constraints. CBQFT-2 inherits the nonlocal features of Bohmian QFT, but its context switching adds a new layer that must be carefully assessed. As noted in Section 5, the nonlocality characteristic of Bohmian mechanics already appears in Bell-type QFT and in CBQFT-1: distant measurement settings can instantaneously alter the jump rates for particle transitions due to entanglement of the global wavefunction. CBQFT-2 deepens this picture by making the *representation* of the field algebra itself context-sensitive. A transition in the macroscopic context variable Λ_t—for instance, corresponding to a phase transition or a measurement interaction—entails a shift from one (typically unitarily inequivalent) representation of the quantum field algebra $\mathcal{A}$ to another.

Because the quantum state Ψ_t is always defined in a context-sensitive Hilbert space $\mathcal{H}_{\Lambda_t}$, a change in Λ_t prompts a re-expression of Ψ_t into the new sector (Axiom B6). At the algebraic level, however, the physical content of Ψ_t is continuous across the jump: for all observables A in the chosen overlap algebra $\mathcal{A}_0$, one has $\omega_{t^-}(A) = \omega_{t^+}(A)$, so local expectations outside the trigger region are unchanged.

One might refer to this as *representation-theoretic nonlocality*. A symmetry-breaking event in one spatial region—triggered, for example, by a particle trajectory entering a detector—can drive a stochastic transition in Λ_t via the context kernel $P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t)$, and this in turn re-labels the active global representation π_{Λ_t}. In that sense the "background representation" is updated globally (relative to the preferred foliation). However, two important caveats ensure compatibility with relativistic causality and foreclose superluminal signalling.

(i) Re-expression at a context jump is implemented by an isometry $U_{\Lambda_i \to \Lambda_f}$ that preserves all expectation values on any local algebra $\mathcal{A}_0$ chosen outside the trigger region. Thus no agent can use a local intervention on Λ_t to change the statistics of local measurements at spacelike separation.
(ii) The contextual transition is a *stochastic* update of the representation label, not a controllable signal. It is driven by microscopic events (such as a particle entering a detector) according to the kernel $W_{\Lambda_i \to \Lambda_j}(Q_t, \mathcal{J}_t)$, and cannot be freely tuned to transmit information faster than light.

Thus no agent can exploit a local intervention on Λ_t in one region to change the outcome statistics of any experiment confined to a spacelike separated region: all local observables there are represented in $\mathcal{A}_0$, and their expectation values are invariant under the context switch. The nonlocality of CBQFT-2 is therefore of the same basic kind as in Bell-type Bohmian QFT: the dynamics of the configuration Q_t and of the context Λ_t are defined on a preferred foliation, and the global state responds instantaneously along that foliation. What CBQFT-2 adds is that the *choice of representation* itself becomes part of the nonlocal structure, while local expectations remain insensitive to distant context switches.

Nevertheless, this feature raises subtle questions about compatibility with relativistic covariance. In the heuristic formulation of CBQFT-2, the context is already described by a spacetime field $\Lambda(x, t)$ encoding detector settings, phases, and other macroscopic structures. In the present formalism we compress this information, at each time *t*, into a single label Λ_t that selects the representation π_{Λ_t} and Hilbert space $\mathcal{H}_{\Lambda_t}$ on the preferred time-slice.

A fully covariant extension, we suggest, would make this dependence on $\Lambda(x, t)$ more explicit and more local: to each spacetime region $\mathcal{O}$ one would assign not only a local algebra $\mathcal{A}(\mathcal{O})$, as in algebraic QFT, but also a context-dependent representation $\pi_{\Lambda|\mathcal{O}} : \mathcal{A}(\mathcal{O}) \to \mathcal{B}(\mathcal{H}_{\Lambda|\mathcal{O}})$, with overlap conditions ensuring that expectations agree on intersections $\mathcal{O}_1 \cap \mathcal{O}_2$.

Within such a framework one can still retain a Bohmian configuration Q_t defined on a preferred foliation: on each leaf Σ_t the global state Ψ_t lives in the representation determined by the spatial profile $\Lambda(\cdot, t)$, and Q_t is guided by Ψ_t in exactly the usual way (with additional Bell-type jumps). The context field $\Lambda(x, t)$ then evolves stochastically according to local rules constrained by the spacetime causal structure, while the Bohmian dynamics of Q_t is defined along the foliation. This parallels the situation in relativistic Bohmian QFT more generally: Lorentz covariance is not manifest at the level of the beables, but no-superluminal-signalling is maintained for all local observables, and the empirical content is compatible with relativistic constraints.

9 Other QFT Interpretations

We have been considering how CBQFT extends standard formulations of Bohmian quantum theory. Besides standard BM and Bell-type BQFTs, however, two prominent families of interpretations have been applied to QFT: the Everett (Many-Worlds) interpretation and spontaneous collapse models such as GRW and CSL. They likewise aim to solve the measurement problem. How do they compare with CBQFT?

9.1 Everett (Many-Worlds)

In Everettian QFT, the universal wavefunction evolves unitarily and encodes all possible outcomes of particle interactions in a vast, continually branching superposition. There are no stochastic transitions, no particle trajectories, and no unique outcomes—just a continually decohering structure in which each branch corresponds to a different macroscopic world [59, 60].

CBQFT-2 diverges sharply from this view. It maintains a *single, actual* particle (or field) configuration Q_t, guided by a state $\Psi_t \in \mathcal{H}_{\Lambda_t}$ (and, across context switches, re-expressed into an inessentially amplified space $\widetilde{\mathcal{H}}_{\Lambda_t}$), where Λ_t is a dynamically evolving context variable. Contexts—unlike branches—are not mere patterns in the state but physically instantiated macroscopic conditions that modulate and are modulated by microphysical processes. The context kernel $P(\Lambda_f \mid \Lambda_i, Q_t, \mathcal{J}_t)$ encodes the probability of context shifts, producing genuine stochastic transitions between (typically) inequivalent representation sectors. Within each sector, evolution is unitary. In the hybrid formulation of Section 7.7, the unconditional evolution on classical–quantum (block-diagonal) states defines a CPTP map on the enlarged register–sector space, rather than a single universal unitary on a fixed Hilbert space.

Everettian accounts face long-standing challenges in deriving the Born rule, explaining the empirical salience of branches, and specifying a preferred decomposition into system and environment [1, 28]. CBQFT-2 does not attempt a new derivation of the Born rule; it inherits the standard Bohmian solution via equivariance and a quantum-equilibrium assumption (Sections 4–7). It *does* avoid Everett's ambiguity about branching structure by treating the system–context decomposition as part of the theory's architecture: macroscopically distinct situations correspond to representation sectors of the field algebra, indexed by Λ_t, with explicit stochastic dynamics.

Furthermore, whereas Everettianism regards macroscopic structures as weakly emergent patterns in the universal wavefunction, CBQFT-2 treats them as being ontologically more robust: contexts Λ_t correspond to macroscopically distinguishable sectors of reality, realised via boundary, thermodynamic, or measurement-related structures. This preserves the causal efficacy and identities of macroscopic systems instead of dissolving them into a proliferation of parallel worlds.

9.2 Collapse Models (GRW/CSL/CWC)

Spontaneous collapse theories resolve the measurement problem by modifying the quantum dynamics. GRW introduces rare, random collapses of the wavefunction for each particle [22], while CSL replaces this with continuous stochastic noise [40]. These approaches yield effective localisation and definite outcomes, but they raise significant questions about the underlying ontology, the status of energy conservation, Lorentz invariance, and so on.

CBQFT-2 offers a sharply contrasting mechanism. Rather than modifying the Schrödinger equation or introducing *ad hoc* noise terms, CBQFT-2 embeds collapse-like behaviour within a rigorous stochastic process over macroscopic contexts. The variable Λ_t evolves via a transition kernel informed by the current configuration Q_t and contextual data $\mathcal{J}_t$. Transitions between Λ_t values induce an isometric re-expression of the guiding state: at a context jump $\Lambda_i \to \Lambda_f$, $\Psi_{t^-} \in \mathcal{H}_{\Lambda_i}$ is mapped to $\Psi_{t^+} \in \widetilde{\mathcal{H}}_{\Lambda_f}$ so as to preserve all overlap expectations on a local algebra $\mathcal{A}_0$ (Axiom B6). This yields effective collapse-like behaviour via stochastic context selection together with exact sector classicality (superselection) in the context label (Section 7.5.1), while preserving norm and full Bohmian equivariance (Sections 7.4.1–7.4.2).

Because only one particle configuration and one context exist at a time, CBQFT-2 sidesteps many of the worries associated with the "tails problem" in GRW/CSL: small-amplitude components of the wavefunction do not correspond to ghostly objects; the actual history is carried by a single Q_t and Λ_t, with the rest of the state playing the usual guiding and probabilistic role familiar from BM. CBQFT-2 maintains a robust, structured ontology grounded in dynamically interacting particles and macroscopic structures.

A closely related framework is Contextual Wavefunction Collapse (CWC) theory [16], which interprets wavefunction collapse as an emergent, thermodynamically driven process resulting from contextual interactions between a quantum system and its macroscopic environment. Like CWC, CBQFT-2 treats collapse as a dynamical response to changes in environmental structure—but it refines this idea by embedding such contextual changes in a rigorous algebraic formalism and by coupling them to a Bohmian configuration. Rather than merely tracing over environmental degrees of freedom, CBQFT-2 implements decoherence and effective collapse through transitions between (typically) unitarily inequivalent Hilbert-space sectors associated with macroscopically distinct contexts. In this way, CBQFT-2 offers a concrete realisation of CWC-like principles within a Bohmian framework.

10 Concluding Remarks

Contextual Bohmian QFTs (CBQFTs) extend Bell-type Bohmian quantum field theory by maintaining their particle realism whilst modifying their stochastic dynamics to admit an explicit macroscopic contextual dependence. We have presented here two ways of formulating this approach.

In CBQFT-1, the macroscopic physical context enters the dynamics as an externally specified parameter Λ that modulates the Hamiltonian and alters particle jump rates. This model preserves the Bohmian ontology while integrating environmental or experimental conditions in a controlled, top-down manner.

CBQFT-2 generalises this picture by treating context Λ_t as a dynamical, piecewise-continuous stochastic process that evolves in response to the quantum configuration. Each context value determines a distinct (generally unitarily inequivalent) representation π_{Λ_t} of the field algebra, defining a Hilbert space $\mathcal{H}_{\Lambda_t}$ in which the guiding state Ψ_t is represented. When Λ_t jumps, Ψ_t is *re-expressed* by a (partial) isometry that preserves all overlap-algebra expectations on a chosen local $\mathcal{A}_0$. This dynamical architecture closes a feedback loop between particle events and macroscopic conditions while retaining unitary evolution between context changes and securing norm conservation and full Bohmian equivariance. Here, contextual influences are treated as part of the theory's internal architecture rather than as an external modelling choice.

Philosophically, CBQFTs preserve the realism of Bohmian mechanics, avoid the extravagance of Everettian many-worlds, and sidestep the *ad hoc* collapse postulates of GRW-type models. They offer a concrete, formally grounded and well-motivated account of how macroscopic features—like measurement apparatus, boundary conditions, or thermal environments—can shape quantum dynamics without introducing observer dependence.

More speculatively, the framework we have proposed in this paper revives a vision of nature in which the physical world consists of matter structured by form [53]: macroscopic form, encoded in the context variable Λ_t, not only arises from but also helps govern the behaviour of quantum matter. Given this shift from context as an external condition to form as a dynamically evolving structural field, it is natural to view CBQFT-2 as more than just an extension of Bohmian mechanics. In Section 8, we recommended classifying it as a *hylomorphic* quantum field theory: that is, a theory in which the forms of macroscopic (middle-sized) systems play an active, lawfully mediated role in shaping the evolution of physical systems. Hylomorphism thus opens the door to a causally bidirectional, thermodynamically informed, and ontologically layered understanding of QFT, in which middle-sized things, such as scientists and their instruments of measurement, make a causal difference to how nature unfolds.

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