

The Field Ontology in Quantum Field Theory

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Contents

A. Particles, Fields or both?	1
A.1. A Provocative Example: Bohm’s 1952 paper	2
A.1.1. Digression: Bosons as Virtual Particles — The Beauty of a Field Ontology?	4
B. The Lagrangian QFT Approach and its Natural Pilot-Wave Theories	7
B.1. Canonical Quantization of Fields (Non-Rigorous)	7
B.2. How Canonical Quantization Shouts for a Pilot-wave Theory	10
B.2.2. Digression: Why is the quantization of fields called “second quantization”?	12
B.3. What about Fermions?	12
C. Rigorous Field-configuration QM I: Mode Expansion Cut-off	14
C.1. Ultraviolet (UV) Cut-off	14
C.1.1. Digression: The Quantum Potential as the Origin of Mass	16
C.2. Discrete Space	16
C.3. QFT and the Endless Tower of Approximate Theories	17
D. Rigorous Field-configuration QM II: Insist on Infinite Degrees of Freedom	18
D.1. Problem I: The Operator Fields	18
D.2. Problem II: The Measure	19
D.2.1 The Key Strategy	21
D.2.2. The Price to Be Paid	22
D.3. Algebraic QFT and How to Use it to Build a Distinguished QFT over Field Configurations	23
D.3.1. The C^* -algebras that Abstract the Commutation Relations	24
D.3.2. A rigorous QFT over Fock-Space	26
D.3.3. A Baby Version of the Grand Finale	27
D.4. A rigorous QFT with configuration-space \mathbb{R}^∞ : “a Ceiling for the Infinite Tower?!?”	30
D.4.1. A Rigorous Pilot-Wave Theory in \mathbb{R}^∞ ?	33
E. A Meditation in the Summit: “there were Higher Peaks Above the Ceiling..”	34
F. Bibliography	36
G. Appendices	38
Appendix to §B.1: Calculus in Infinite Dimensional Vector Spaces	38
Appendix to §B.2: Not all PWTs explain the collapse and “measurement” probabilities!	41
Appendix to §D.2: How the Trick is Employed in the Feynman-Kac Formula	44

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“Scientific Project” — directed by Prof. Roderich Tumulka

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A. PARTICLES, FIELDS OR BOTH?

The so-called *Bohmian* theories (Tumulka, 2022; Dürr et al., 2005, 2004; Deckert et al., 2016) are able to explain all from non-relativistic quantum mechanics (QM) till quantum field theory (when equipped with a suitable configuration POVM). For that, their key postulate is that there are actual point-like particles that exist irrespective of our aims to detect them, explaining in doing so, the visible image of the world and the results of quantum experiments.¹ So, the usual choice in Bohmian theories is a point-particle primitive ontology.

But, could we consider that instead of infinitesimally small entities what *there is* are actually just the opposite: infinitely extended entities filling the whole space? In a classical setting such entities were commonplace: the *physical-space fields* (as opposed to configuration-space fields), like the electric and magnetic field or a density field in fluid mechanics. In the present manuscript, instead of particle configuration-space, such as \mathbb{R}^{3N} (where each point represents a particular position for N particles in physical space), we will consider *field configuration-space*, say $L^2(\mathbb{R}^3, \mathbb{R})$, a space where each point gives a field that occupies the whole physical space. Instead of a “superposition” of particle configurations (represented by a wavefunction giving a complex number for each configuration), we will have a “superposition” of fields (given by a wavefunction-*al* assigning a complex number to each field). As such, these objects will be very natural to Quantum *Field* Theory (QFT).

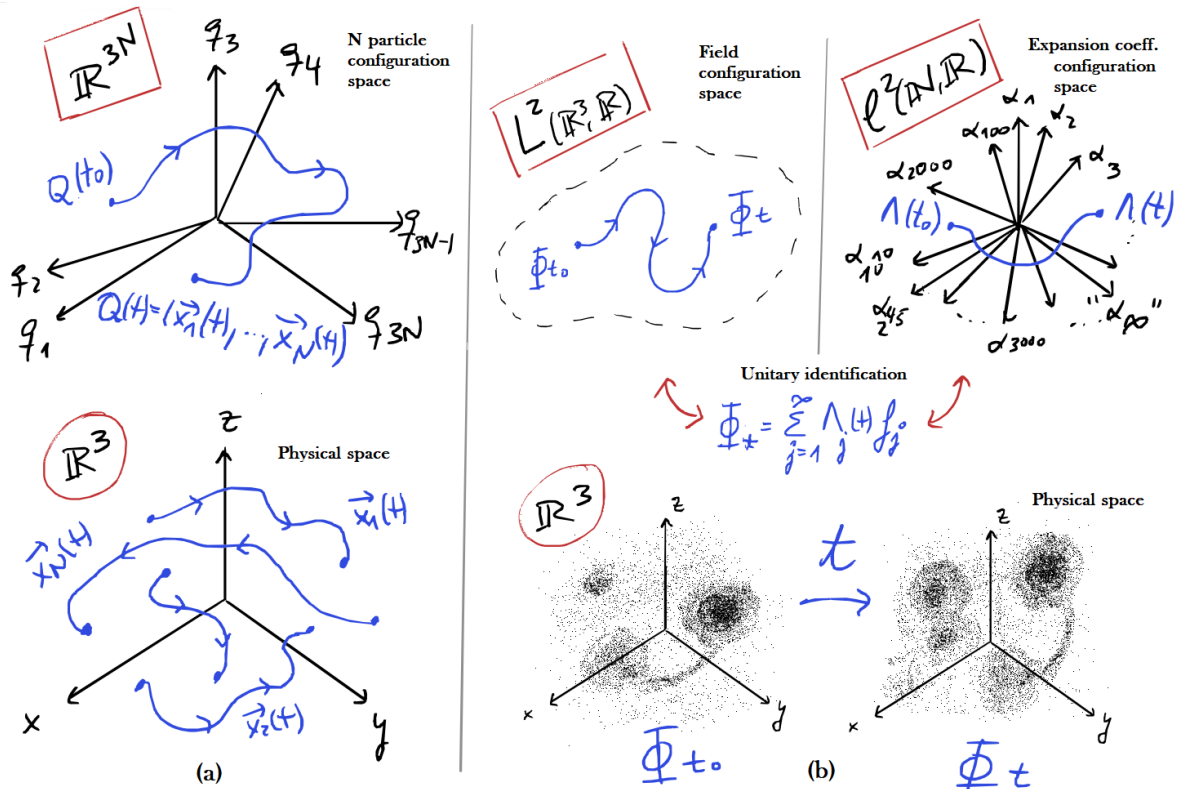


Figure 1: (a) shows a representation of configuration-space of N point-like particles, given by \mathbb{R}^{3N} . A point there represents a particular arrangement of the N positions of the particles in physical-space. (b) shows a representation of a field configuration-space $L^2(\mathbb{R}^3, \mathbb{R})$ and its associated expansion-coefficient configuration-space $\ell^2(\mathbb{N}, \mathbb{R})$ for some ONB $\{f_j\}_{j \in \mathbb{N}} \subset L^2(\mathbb{R}^3, \mathbb{R})$. Each point in either space corresponds to a whole physical-space field via the unitary identification written in the figure.

¹The main merit of Bohmian mechanics is that it explains mechanistically the collapse postulate of textbook QM (Dürr et al., 1992; Tumulka, 2022), solving the measurement problem just as a simple corollary of the theory.

A.1. A Provocative Example: Bohm's 1952 paper

In an appendix of his seminal paper of two parts, [Bohm \(1952\)](#) already gave us a taste of what a “field ontology” would mean in QM, in the very same article presenting what a “particle ontology” means. Let us describe it as an introduction, for it will serve as an example of the topics covered throughout the document.

Consider the classical description of N point-like particles (say, electrons) of position-velocity coordinates $\{(\vec{r}_j, \dot{\vec{r}}_j)\} \subset \mathbb{R}^{6N}$ and electric charges $\{e_j\}_j$ and masses $\{m_j\}_j$. Consider also an electromagnetic 4-potential field $(\phi_{\vec{r}}, \vec{A}_{\vec{r}}) : \mathbb{R}^3 \rightarrow \mathbb{R}^4$, composed by fields on some Sobolev space. The Lagrangian of such a classical system is given by

$$L(\vec{r}_j, \dot{\vec{r}}_j, \vec{A}_{\vec{r}}, \dot{\vec{A}}_{\vec{r}}, \phi_{\vec{r}}) = \sum_{j=1}^N \left(\frac{1}{2} m_j |\dot{\vec{r}}_j|^2 + \frac{e_j}{c} \dot{\vec{r}}_j \cdot \vec{A}_{\vec{r}_j} - e_j \phi_{\vec{r}_j} \right) + \frac{\varepsilon_0}{2} \int (|\vec{E}_{\vec{r}}|^2 - c^2 |\vec{B}_{\vec{r}}|^2) d^3r \quad (1)$$

where the electric and magnetic field are defined respectively as $\vec{E}_{\vec{r}} := -\dot{\vec{A}}_{\vec{r}} - \vec{\nabla} \phi_{\vec{r}}$ and $\vec{B}_{\vec{r}} := \vec{\nabla} \times \vec{A}_{\vec{r}}$.² The Euler-Lagrange equations³ are exactly the Maxwell equations (the two that do not follow from the definitions of \vec{E} and \vec{B}) and the Newton's second Law with the Lorentz and Coulomb forces in it. The resulting equations are poorly defined because each particle is centred in a singularity of the total force field (among other things). Let us ignore that for a moment. Fix the Coulomb gauge ($\vec{\nabla} \cdot \vec{A} = 0$) and explicitly write $\phi_{\vec{r}} = \sum_k \frac{e_k}{4\pi\varepsilon_0 |\vec{r} - \vec{r}_k|}$, which is the solution to the Gauss equation (i.e., the Poisson equation in Coulomb gauge) for a point-charge distribution $\{\vec{r}_j\}_j$. This way we get rid of ϕ as a degree of freedom. Then we can get rid of some of the problems of electrodynamics by assuming there is no “self-interaction” by “re-scaling” the energy (the first historic “renormalization”) so that defining $V_{j,k}(|\vec{r}_j - \vec{r}_k|) := e_j e_k / (4\pi\varepsilon_0 |\vec{r}_j - \vec{r}_k|)$, the evaluation of ϕ in the Lagrangian turns

$$\text{from } \frac{1}{2} \sum_j \sum_k V_{j,k} \quad (\text{infinite whenever } j = k) \text{ to } \frac{1}{2} \sum_j \sum_{k \neq j} V_{j,k}. \quad (2)$$

Now, as explained in mathematical rigour for instance in [\(Abraham et al., 2012\)](#), one can consider a symplectic structure for the infinite dimensional phase space of $\vec{A}_{\vec{r}}$ and its canonically conjugate variable $\vec{\pi}_{\vec{r}} := \dot{\vec{A}}_{\vec{r}}$, together with the usual phase space for the particles. A Legendre transformation leaves us a Hamiltonian that after some manipulations (integration by parts) looks like:

$$H = \sum_{j=1}^N \frac{(\vec{p}_j - e_j \vec{A}_{\vec{r}_j})^2}{2m_j} + \frac{1}{2} \sum_j \sum_{k \neq j} V_{j,k} (|\vec{r}_j - \vec{r}_k|) + \frac{\varepsilon_0}{2} \int (\dot{\vec{A}}_{\vec{r}} + c^2 |\vec{B}_{\vec{r}}|^2) d^3r \quad (3)$$

Assume we want to describe electrodynamics in a cavity (or in a finite Universe...), so that physical-space can be considered instead of \mathbb{R}^3 to be $\Omega := (-\frac{L}{2}, \frac{L}{2})^3$ for some fixed $L > 0$. Then, we can artificially consider periodic boundary conditions so that using only Fourier modes of commensurate wavelengths with the box, $\vec{k} \in (\frac{2\pi}{L})\mathbb{Z}^3$, we can expand in Fourier series any L^2 function inside the box. In this line, from PDE theory one gets that there is a countable number of solutions indexed by $\vec{k} \in \mathbb{Z}^3 (2\pi/L)$ for the constrained eigenvalue problem

$$\begin{cases} \Delta \vec{u}_{\vec{k}}(\vec{r}) = -|\vec{k}|^2 \vec{u}_{\vec{k}}(\vec{r}) \\ \vec{\nabla} \cdot \vec{u}_{\vec{k}}(\vec{r}) = 0 \end{cases}, \quad (4)$$

yielding an orthonormal basis (ONB) of $\{\vec{f} \in L^2(\Omega, \mathbb{R}^3) \mid \vec{\nabla} \cdot \vec{f} = 0\}$. In fact, although they are vector valued, they will be made of complex exponentials (or rather sines and cosines) —so-called plane-waves or Fourier modes. Because they form an ONB of those $\vec{A}_{\vec{r}} \in L^2(\Omega, \mathbb{R}^3)$ that satisfy the Coulomb gauge, there is always a unique sequence of coefficients $\{\alpha_{\vec{k}}\}_{\vec{k} \in (\frac{2\pi}{L})\mathbb{Z}^3} \subset \mathbb{C}$ so that

$$\vec{A}_{\vec{r}} \equiv \sum_{\vec{k} \in (\frac{2\pi}{L})\mathbb{Z}^3} \alpha_{\vec{k}} \vec{u}_{\vec{k}}(\vec{r}) \implies \vec{A}_{\vec{r}}(t) \equiv \sum_{\vec{k} \in (\frac{2\pi}{L})\mathbb{Z}^3} \alpha_{\vec{k}}(t) \vec{u}_{\vec{k}}(\vec{r}) \quad \left\{ \begin{array}{l} \text{If the field depends on time, the depen-} \\ \text{dence will be “charged” to the coefficients!} \end{array} \right. \quad (5)$$

²Two of the Maxwell equations are engraved in these definitions. By vector calculus, $\vec{\nabla} \times \vec{E}_{\vec{r}}(t) \equiv -\frac{d}{dt} \vec{B}_{\vec{r}}(t)$ and $\vec{\nabla} \cdot \vec{B}_{\vec{r}}(t) \equiv 0$.

³We will see in a moment how to take derivatives with respect to field variables rigorously.

Since there is a one-to-one correspondence between $(\alpha_{\vec{k}})_{\vec{k}} \in \ell^2\left(\frac{2\pi}{L}\mathbb{Z}^3, \mathbb{C}\right)$ and $\vec{A}_{\vec{r}} \in L^2(\Omega, \mathbb{R}^3)$, we can equally use the expansion coefficients as the degrees of freedom to specify a field-configuration instead of the vector field \vec{A} over \mathbb{R}^3 . Plugging this in the Hamiltonian (3), we get a Hamiltonian now as a function of the expansion coefficients:

$$H(\vec{r}_j, \alpha_k, \vec{p}_j, \alpha_k) = \sum_{j=1}^N \frac{\left(p_j - e_k \overbrace{\sum_k \alpha_{\vec{k}} \vec{u}_{\vec{k}}(\vec{r}_j)}^{\vec{A}_{\vec{r}_j}}\right)^2}{2m_j} + \frac{1}{2} \sum_j \sum_{k \neq j} V_{j,k}(|\vec{r}_j - \vec{r}_k|) + \frac{1}{2} \sum_{\vec{k} \in \frac{2\pi}{L}\mathbb{Z}^3} \left(\frac{\pi^2}{k^2} + c^2|\vec{k}|^2 \alpha_{\vec{k}}^2\right). \quad (6)$$

(One can find that what we did is essentially a canonical transformation in infinite dimensional vector space.) Notice that the Hamiltonian is describing a harmonic oscillator for each of the mode coefficients (coupled to the other degrees of freedom). Since, we know a working quantization prescription in such a case, we can get the quantum counterpart of the theory. But first, by reasons we will see later (namely, that there is no Lebesgue measure in infinite dimensions), we must assume what is called an ‘‘ultraviolet (UV) cut-off’’: essentially that all coefficients above certain finite $K > 0$ (in terms of the norm of \vec{k}) are roughly zero, so we are left with only M degrees of freedom that fully describe the vector field \vec{A} , say $\{\alpha_k\}_{k=1}^M$. Equivalently, since this puts a limit on the frequency of the Fourier modes we are using, it is a limit on the spatial variation of the vector \vec{A} . If this makes one feel discomfort (as it should), assume that K is very big, say, $K = 10^{10^{100}}$, so that the limitation is posed only for length-scales that are absurdly small.

The cut-off leaves us with finitely many degrees of freedom, so we can perform a canonical quantization *à la* Dirac,⁴ and we will get (not-fully-relativistic) quantum electrodynamics (QED). That is, the quantum counterpart of classical electrodynamics. It will have each system described by a wavefunction $\psi_t \in L^2(\mathbb{R}^{3N+M}, \mathbb{C}, d^{3N+M}x)$ ($d^n x$ is the Lebesgue measure for \mathbb{R}^n), with arguments $\psi_t(\vec{r}_1, \dots, \vec{r}_j, \alpha_1, \dots, \alpha_M)$ that for all $t \in \mathbb{R}$ obeys the following Schrödinger Equation (SE):

$$i\hbar \frac{d}{dt} \psi_t(\{\vec{r}_j, \alpha_k\}) = \left[\sum_{j=1}^N \frac{\left(-i\hbar \vec{\nabla}_{\vec{r}_j} - e_j \vec{A}(\vec{r}_j; \{\alpha_k\}_k)\right)^2}{2m_j} + \frac{1}{2} \sum_j \sum_{k \neq j} V_{j,k}(|\vec{r}_j - \vec{r}_k|) + \sum_{k=1}^M \left(-\frac{\hbar^2}{2} \frac{\partial^2}{\partial \alpha_k^2} + \frac{c^2|\vec{k}|^2}{2} \alpha_k^2\right) \right] \psi_t(\{\vec{r}_j, \alpha_k\}) \quad (7)$$

with $\vec{A}(\vec{r}_j; \{\alpha_k\}_k) := \sum_{k=1}^M \alpha_{\vec{k}} \vec{u}_{\vec{k}}(\vec{r}_j)$.

Now, one could wonder how robust this is upon a change of the expansion basis chosen before quantization, or under a change in the gauge condition, or a change of the size L of the cavity. And those are very good questions that do matter, but since the SE we found is exactly what physicists use in solid state and quantum optics (Cohen-Tannoudji et al., 1989), perhaps our *Ansätze* were somewhat acceptable.

Back to our point: since we have a finite number of degrees of freedom, this is a ‘‘conventional’’ quantum theory, so that we can apply the idea of Bohmian mechanics. That is, we find there is a continuity equation for $|\psi_t(\{\vec{r}_j, \alpha_k\})|^2$ (obtained by taking $\frac{d\psi_t \psi_t^*}{dt} = \psi_t \frac{d\psi_t^*}{dt} + \psi_t^* \frac{\partial \psi_t}{\partial t}$ and plugging the SE there)

$$\frac{d|\psi_t|^2}{dt} = -\vec{\nabla} \cdot (|\psi_t|^2 \vec{v}) \quad \text{where } \vec{v} := (\vec{v}^{r_1}, \dots, \vec{v}^{r_N}, v^{\alpha_1}, \dots, v^{\alpha_M}) \quad (8)$$

$$\vec{v}^{r_i}(\{\vec{r}_j, \alpha_k\}) := \text{Re} \left\{ \frac{-i\hbar \vec{\nabla}_{\vec{r}_i} \psi_t(\{\vec{r}_j, \alpha_k\})}{m_j \psi_t(\{\vec{r}_j, \alpha_k\})} \right\} \quad \text{and} \quad v^{\alpha_l}(\{\vec{r}_j, \alpha_k\}) := \text{Re} \left\{ \frac{-i\hbar \frac{\partial}{\partial \alpha_l} \psi_t(\{\vec{r}_j, \alpha_k\})}{\psi_t(\{\vec{r}_j, \alpha_k\})} \right\}. \quad (9)$$

Next, we postulate that there is an actual, determined set of N point-particles and an actual electromagnetic field, of joint trajectory $t \mapsto \left(\vec{R}_j^\xi(t), \Lambda_k^\xi(t)\right)_{j,k}$ following the velocity field that guides the density in the continuity equation, i.e.,

$$\begin{cases} \frac{d\vec{R}_j^\xi(t)}{dt} = \vec{v}^{r_j} \left(\left\{ \vec{r}_j = \vec{R}_j^\xi(t), \alpha_j = \Lambda_k^\xi(t) \right\} \right) \\ \frac{d\Lambda_k^\xi(t)}{dt} = v^{\alpha_j} \left(\left\{ \vec{r}_j = \vec{R}_j^\xi(t), \alpha_j = \Lambda_k^\xi(t) \right\} \right) \\ (\vec{R}_j^\xi(t_0), \Lambda_k^\xi(t_0)) = \vec{\xi} \end{cases} \quad \text{for } j \in \{1, \dots, N\}, k \in \{1, \dots, M\} \quad (10)$$

⁴That is (satirically), ‘‘hocus pocus, now positions and momenta are operators satisfying certain commutation relations’’.

for some “initial condition” $\vec{\xi} \in \mathbb{R}^{3M+N}$. Then, $\vec{A}_{\vec{r}}^{\xi}(t) := \sum_k \Lambda_k^{\xi}(t) \vec{u}_k(\vec{r})$ would be the actual vector-potential field selected by the trajectory of the expansion coefficients $\Lambda_k^{\xi}(t)$ (and hence there would be a particular electric and magnetic field \vec{E}, \vec{B} selected by this $\vec{A}_{\vec{r}}(t)$). Because the trajectory is a flow-line of the velocity field guiding the density, if $\vec{\xi}$ was sampled following the probability measure $|\psi_{t_0}|^2 d^{3N+M}x$, the trajectory $\left(\vec{R}_t^{\xi}(t), \Lambda_k^{\xi}(t)\right)_{j,k}$ will be distributed according to $|\psi_t|^2 d^{3N+M}x$ at all times t . Finally, assuming that measurement device pointers are made of some of the N particles, the measurement problem and the collapse postulate would be a corollary just as in Bohmian mechanics.

A.1.1. Digression: Bosons as Virtual Particles — The Beauty of a Field Ontology?

All in all, we have found that in the presented theory, the charged particles have a particle ontology (i.e., the particles exist as points at all times), while the electromagnetic field has a field ontology (i.e., the electric and magnetic fields exist as physical-space fields at all times). Akin to the classical theory. This is opposed to what one thinks about when we talk about a *photon*, a light particle, often said to be what emerges when we quantize the electromagnetic field. But where is it here? In this ontology a photon is nothing but an apparent particle, a *virtual particle*.

Recall that $L^2(\mathbb{R}^{3N+M}) = L^2(\mathbb{R}^{3N}) \otimes L^2(\mathbb{R}^M)$, with $L^2(\mathbb{R}^M)$ the Hilbert space over the field configurations—or more precisely their expansion coefficients. Now, from the Hamiltonian in (7), we see that each coefficient α_k behaves as a harmonic oscillator coupled to the particle velocities. If there were no charged particles, $N = 0$, we would have the electromagnetic (EM) field alone, and the α_k would behave as independent harmonic oscillators with a full Hamiltonian $H(\{\alpha_k, \pi_k\}) = \sum_k H_{free\ k}$, where $H_{free\ k} := \frac{1}{2}(\pi_k^2 + c^2|k|^2\alpha_k^2)$. The latter has an ONB of eigenfunctions on $L^2(\mathbb{R})$ (which is the k -th mode’s Hilbert space), that we denote by $\{\phi_n^k\}_{j \in \mathbb{N}_0} \subset L^2(\mathbb{R})$ (see Figure 2.a). They are well-known to be Gaussians times Hermite polynomials $H_n(x)$:

$$\phi_n^k(\alpha_k) := \frac{1}{\sqrt{2^n n! \pi^{1/4}}} e^{-\frac{1}{2}\alpha_k^2} H_n(\alpha_k) = \frac{1}{\sqrt{n! 2^n \pi^{1/4}}} \left(\alpha_k - \frac{\partial}{\partial x}\right)^n e^{-\alpha_k^2/2}. \quad (11)$$

With that, we get an ONB now in the full Hilbert space of the field coefficients $L^2(\mathbb{R}^M)$ by taking tensor products of the possible combinations of “single-mode ONBs”: $\{\phi_{j_1}^{k_1} \otimes \dots \otimes \phi_{j_M}^{k_M}\}_{j_1, \dots, j_M \in \mathbb{N}} \subset L^2(\mathbb{R}^M)$, so that any wavefunction over the field-configurations can be expanded as a superposition of them. Now, what physicists call “electromagnetic vacuum” is nothing but the ground state of a free electromagnetic field (i.e., where there are no charges) as explained in (Cohen-Tannoudji et al., 1989), which given by the wavefunction

$$\psi(\alpha_{k_1}, \dots, \alpha_{k_N}) = \phi_0^{k_1}(\alpha_{k_1}) \dots \phi_0^{k_M}(\alpha_{k_M}) = C e^{-\sum_k \alpha_k^2/2}. \quad (12)$$

Because its associated probability density $|\psi|^2$ is a Gaussian around 0, most Bohmian trajectories of the modes $\{\alpha_k\}_k$ will be close to zero $\Lambda_k^{\xi} \simeq 0$, implying that typical Bohmian worlds will have an electromagnetic potential $\vec{A}_{\vec{r}} \simeq 0$ (and hence electric and magnetic fields $\vec{E}, \vec{B} \simeq 0$). But there are a significant part of trajectories Λ_k^{ξ} that are a bit further from the origin, which causes the so-called “vacuum fluctuations”.

On the other hand, we say there is only *one* photon and has frequency \vec{k}_j if the wavefunction is

$$\psi(\alpha_{k_1}, \dots, \alpha_{k_N}) = \phi_0^{k_1}(\alpha_{k_1}) \dots \phi_1^{k_j}(\alpha_{k_j}) \dots \phi_0^{k_M}(\alpha_{k_M}) \propto \alpha_k e^{-\sum_k \alpha_k^2/2}. \quad (13)$$

two photons of frequency \vec{k}_j if

$$\psi(\alpha_{k_1}, \dots, \alpha_{k_N}) = \phi_0^{k_1}(\alpha_{k_1}) \dots \phi_2^{k_j}(\alpha_{k_j}) \dots \phi_0^{k_M}(\alpha_{k_M}) \propto (\alpha_k^2 - 1) e^{-\sum_k \alpha_k^2/2}. \quad (14)$$

and so on (see Figure 2.b). That is, a higher excitation of the harmonic oscillator represented by the mode α_k corresponds to an additional photon of that k -th frequency. This is closely related to the Fock space representation as we will see later.

But where is in all this description the photon that is absorbed and emitted in discrete energy packets? That this happens in packets is an experimental fact! Well, let us put back one of the particles in the description. Assume it is an electron in a hydrogen atom, so that it has well-defined energy eigenstates: we denote them by $\varphi_0, \varphi_1, \dots$ starting from the ground state (the s -orbital). Also, assume we have sent against the atom a pulse of a laser

and then enclosed it between mirrors. The laser has a very sharp frequency \vec{k} , so that we have enough leaving only one of the expansion modes for simplicity, say, the α_k (hence, we only consider fields like $\vec{A}_{\vec{r}} \simeq \alpha_k \vec{u}_{\vec{k}}$). Then, one can find that if we start an experiment with a joint wavefunction $\psi(\vec{r}, \alpha_k) = \varphi_0(\vec{r})\phi_1^k(\alpha_k)$, that is, a tensor product of a “single photon” state ϕ_1^k and the ground state of the electron φ_0 , the Schrödinger equation will make them evolve very close to $\varphi_1(\vec{r})\phi_0^k(\alpha_k)$ —hence, a state describing no photon but the electron now in the excited state— and then back. This are the so-called Rabi oscillations (Cohen-Tannoudji et al., 1989), and it is just the atom absorbing and emitting a photon over and over again (it is in between mirrors). This oscillation only happens in a meaningful way when the eigen-energy difference between ground state and first excited state of the electron and that of the mode match. But because the latter is exactly given by the frequency \vec{k} of the mode (constants up or down), the Rabi oscillation essentially only happens if the laser with which we have lighted the atom contains a frequency matching the energy interval for the electron’s orbitals (light-matter resonance). If we left the atom in an open environment and not between mirrors, the photon would be absorbed, then emitted and that explains the spectral lines being so specific. Likewise one can explain the photoelectric effect, among others.

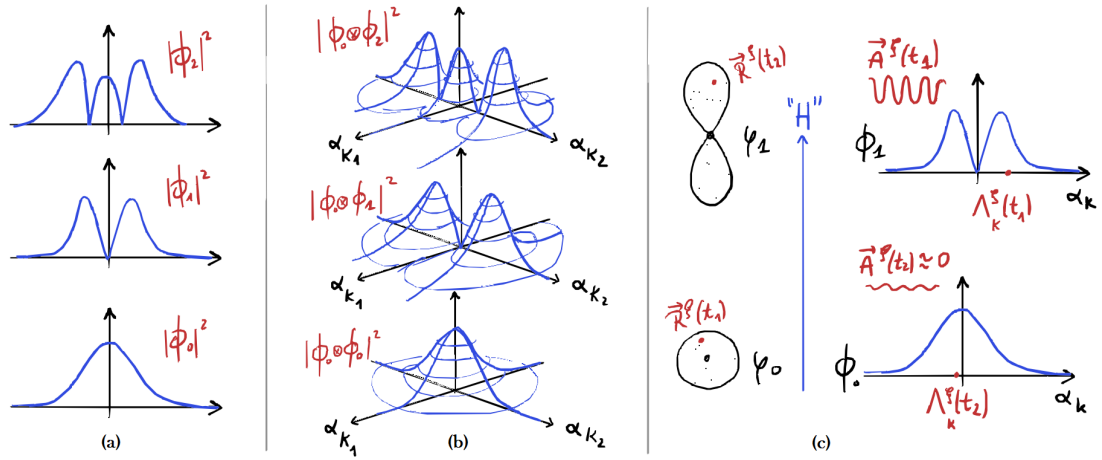


Figure 2: In (a) the magnitude squared of the first three eigenstates of the harmonic oscillator are shown. In (b), for a cut-off at $M = 2$ (two EM field modes), from bottom to top we see: the EM “vacuum” $\phi_0 \otimes \phi_0$; a photon of frequency \vec{k}_2 and none of \vec{k}_1 , i.e., $\phi_0 \otimes \phi_1$; two photons of frequency \vec{k}_2 and none of \vec{k}_1 , $\phi_0 \otimes \phi_2$. In (c) we see the ground state and first excited state of an electron in an atom, next to those of a single mode (the k -th one) of the EM field. The spacing between levels suggests how the eigen-energy difference in one and the other system matches for this particular \vec{k} . We represent example Bohmian positions of the electron and the EM mode coefficient, together with its respective vector potential (if there is a single mode, the vector potential will be a plane wave of the given frequency).

Now, in all the explanation, note there was never an abrupt absorption of anything, all the dynamics had a “soft” time evolution. Moreover, in general, whenever we have an entangled wavefunction between the electron and the modes of the EM field, there is no way to attribute a wavefunction to the electron and photon separately (unless we use conditional wavefunctions or reduced density matrices, which is what one should do to talk about absorption and emission in the above picture more meaningfully).

The point is that in Bohmian terms, with an actual ontological EM field and an actual electron, there is never a discrete particle that is absorbed or emitted! Instead, whenever the wavefunction’s part attributed to the field evolves to a no-photon state (i.e., the joint wavefunction factorizes as $\phi_0^k \varphi_1$), then the ontological field $\vec{A}_{\vec{r}}$ will likely be driven towards zero field strength (the Gaussian is centred in zero), see Figure 2.b. Then, when the system goes back to a single photon state (i.e., the joint wavefunction factorizes as $\phi_1^k \varphi_0$), because the first excited state of the EM mode has two bumps away from zero (see Figure 2), the ontological field $\vec{A}_{\vec{r}}$ is driven to some non-zero field strength. As we see, photons would then be only a “way of talking”. What there would be, is an interaction between the ontological electrons $\{\vec{R}_j(t)\}_j$ and the EM field given by $\{\Lambda_j(t)\}_j$, mediated by the joint wavefunction $\psi(\{\vec{r}_j\}, \{\alpha_k\}_k)$. In terms of ontology, it would look quite similar to what one would expect classically. In fact more than one may think *a priori*!

More specifically, given a solution to the SE (7), one can decompose the wavefunction in polar form by introducing $\psi(\{\vec{r}_j\}, \{\alpha_k\}_k) = R(\{\vec{r}_j\}, \{\alpha_k\}_k) e^{iS(\{\vec{r}_j\}, \{\alpha_k\}_k)/\hbar}$. Plugging this form in the SE and taking the real and imaginary parts of the resulting equation, we get two real partial differential equations (PDEs) obeyed

by $S(\{\vec{r}_j\}, \{\alpha_k\}_k)$ and $R(\{\vec{r}_j\}, \{\alpha_k\}_k)$: one is the continuity equation we mentioned with velocity field \vec{v} and density $R^2 = |\psi|^2$, while the other one is a Hamilton-Jacobi equation for the classical Hamiltonian (6) with an extra “potential energy” term, the so-called Quantum potential, Q^ψ (Holland, 1995). We can evaluate this Hamilton-Jacobi equation along the flow-lines of the Bohmian velocity field \vec{v} , as if we took the Lagrangian frame of fluid mechanics (for a fluid of Bohmian trajectories). We then find an ODE system obeyed along each Bohmian trajectory, which is what inhabitants of a particular Bohmian world could naively⁵ think to be the fundamental equations ruling the particles and the field. These ODEs turns out to be the Newton-Lorentz equation and the Maxwell equations, just that there is the additional force term for the particles due to the “quantum potential” and an additional inhomogeneity for one⁶ of the Maxwell equations.⁷ Namely,

$$m_j \frac{d^2 \vec{R}_j^\xi(t)}{dt^2} = e_j \left(\vec{E}^\xi(\vec{R}_j^\xi(t), t) + \frac{d\vec{R}_j^\xi(t)}{dt} \times \vec{B}^\xi(\vec{R}_j^\xi(t), t) \right) - \vec{\nabla}_{\vec{r}_j} Q^\psi|_{\vec{R}_j^\xi(t), \Lambda_k^\xi(t)} \quad (15)$$

$$\vec{\nabla} \times \vec{B}_{\vec{r}}^\xi(t) = \frac{1}{c^2} \frac{\partial \vec{E}_{\vec{r}}^\xi(t)}{\partial t} + \mu_0 \vec{J}_{\vec{r}}^\xi(t) - F(Q^\psi)|_{\vec{R}_j^\xi(t), \Lambda_k^\xi(t)} \quad (16)$$

where $\vec{J}_{\vec{r}}^\xi := \sum_j e_j \delta(\vec{r} - \vec{R}_j^\xi(t))$ is the usual charge current of classical electrodynamics (with δ the Dirac delta “function”). Q^ψ is the quantum potential and F a map related to the inverse Fourier transform, neither of which we need to detail here.

All in all, if one is more comfortable with it, photons can indeed be considered to be a convenient fiction to describe the interaction of quantized light with electronic orbital shells, but nothing more than that in ontological terms.

But somebody could insist: “what about photon detection screens? I can see photons to be detected as points there!” Well, technically what we really see there is the light emitted by an atom in the screen. The story is that the input EM field’s wavefunction is initially factorized from that of the electron in that atom, but they get entangled when the coupling starts as the laser beam reaches the screen. This generates an entangled wavefunction that performs a Rabi oscillation as explained above, making the atom re-emit the light now in all directions. Since the emission is sourced from essentially a point (the atom), it looks like a point recording to us. But again, light was “absorbed and emitted by the atom” in a non-abrupt manner, and the ontological picture would be close to our classical image of electrodynamics.

Just to be fair, note that an alternative view can also be developed using a Fock space representation of this QFT. Employing the Bell-type QFT technique (Dürr et al., 2005), if we find a position POVM for the photons, we will be able to give them a particle ontology. In such a theory, there would be actual (random) photon creation and annihilation events that look more like literal absorptions and emissions. As of today, there is a problem with finding a position operator for photons (Tumulka, 2022), but still, this might be sooner or later overcome. So at the end of the day it might be up to one’s taste: fields or particles? But importantly, one of them should be fixed in every model. One could try to consider that the same thing (e.g., a photon) can be two things at the same time: particle and field, embracing the confusing orthodox metaphor of “wave-particle duality”. However, if there exist remarkably simpler explanations in which one thing is just one thing, say a particle, or say a field, Occam’s razor must be applied.

⁵This would be a naive believe because given a pair (S, R) solving the continuity and the quantum Hamilton-Jacobi equations, they do not necessarily make up a wavefunction that solves the SE! Hence, it is a nice way to find the classical limit of a quantum theory, because it shows the law of motion of the Bohmian trajectory as seen from within the trajectory, which turns out to have only one extra term with respect to the classical counterpart, but this is not an equivalent PDE system to the SE nor they give an equivalent formulation of QM —in which there is no explicit wavefunction.

⁶Two of the Maxwell equations are readily given as in footnote 1. The Gauss equation was “hardcoded” in our explicit solution to the scalar potential $\phi_{\vec{r}}^\xi(t) = \sum_k \frac{e_k}{4\pi\epsilon_0|\vec{r} - \vec{R}_k^\xi(t)|}$. Hence, only the Ampère-Maxwell law has a quantum “perturbation”.

⁷Both extra terms depend on the wavefunction ψ governing them all across configuration-space, which makes them highly non-local. This breaks for instance the locality properties of the Maxwell equations.

B. THE LAGRANGIAN QFT APPROACH AND ITS NATURAL PILOT-WAVE THEORIES

We will now see in a more general setting that if we start from a field theoretic Lagrangian, i.e., depending on some field-configuration space, say, $L^2(\mathbb{R}^3, \mathbb{R})$, quantization naturally seems to yield wavefunctions $\Psi_t(\phi)$ over such possible field-configurations $\phi \in L^2(\mathbb{R}^3, \mathbb{R})$. In the example above we used a trick to make the degrees of freedom of the field finite, but a priori there is nothing wrong with a complex valued function with arguments in infinite dimensional vector spaces, say $\Psi_t(\alpha_1, \alpha_2, \dots)$ or $\Psi_t(\vec{A}_r)$. These are the so-called *wavefunctionals*. In order to demystify calculus in infinite dimensions you can now proceed to Appendix B.1. There, we rigorously find what physicists call “functional derivatives” and we find a rigorous (and rather simple) derivation of the field theoretic Euler-Lagrange equations.

B.1. Canonical Quantization of Fields (Non-Rigorous)

In this section we will derive the heuristic generalization of canonical quantization to Hamiltonian systems in infinite dimensional symplectic manifolds (the one we exemplified in the introduction). This approach to build QFTs is also called Lagrangian QFT (Wallace, 2006) —although one could rather call it Hamiltonian QFT. Any lecture on QFT by physicists essentially approaches the topic this way, perhaps because it motivates the “field” part of quantum *field* theory, and perhaps because the whole *Standard Model* is built essentially in an analogous way. Nevertheless, as of today, this is not a rigorous approach to QFT (although rigorizations in “tangent” directions do exist, such as the ones we will mention later). We chose to start here because it is the approach to QFT where the field ontology emerges most naturally, by simply generalizing Bohmian mechanics in the line of what we did in the introduction.

Assume we have a finite or infinite dimensional separable Hilbert space⁸ $(Q, \langle \cdot, \cdot \rangle)$, where the points $\eta \in Q$ represent the configuration of some physical quantity at a freezed time. We will call it **configuration-space**. As driving examples we will take $(\mathbb{R}^N, \langle \cdot, \cdot \rangle_2)$ as the configuration space of N particles and $(H^2(\mathbb{R}^3, \mathbb{R}), \langle \cdot, \cdot \rangle_{H^2})$ (the second *Sobolev space*⁹) as the charge density fields. The time evolution of the physical quantity η is assumed then to be given by a path $(t_0, t_f) \subset \mathbb{R} \rightarrow Q; t \mapsto \eta(t)$ in the space¹⁰ Q and will usually be given as an ODE, constituting what we call a “law of physics for η ”. For the example configuration-spaces, we will be interested on Newton’s laws for the particles, and the Klein-Gordon (KG) equation for the charge density (or the wave equation if $m = 0$),

$$m_j \ddot{q}_j(t) = -\frac{\partial V(q)}{\partial q_j} \Big|_{q(t)} \quad \text{and} \quad \ddot{\phi}(t) = \Delta \phi(t) - m^2 \phi(t). \quad (17)$$

Now assume we wish to find candidates to quantum theories whose “classical” counterparts are described by these ODE systems: one says we want to “quantize” what q or ϕ represent. The trick of canonical quantization goes as follows:

1. Find the Lagrangian/action principle of whom, the ODE of the “law of physics” is its Euler-Lagrange equation. In the guiding examples, respectively for $q \in \mathcal{C}^2((t_0, t_f), \mathbb{R}^n)$ and $\phi \in \mathcal{C}^2((t_0, t_f), H^2(\mathbb{R}^3, \mathbb{R}))$, such Lagrangians read

$$L(q, \dot{q}) = \sum_j \frac{1}{2} m_j \dot{q}_j^2 - V(q) \quad \text{and} \quad L(\phi, \dot{\phi}) = \int_{\mathbb{R}^3} \mathcal{L} \left(\phi_{\vec{x}}(t), \dot{\phi}_{\vec{x}}(t), \frac{\partial}{\partial x_j} \phi_{\vec{x}}(t) \right) d^3x \quad (18)$$

would do it, with

$$\mathcal{L}(\phi, \dot{\phi}, \partial_j \phi) = \frac{1}{2} \left(\dot{\phi}^2 - \sum_j (\partial_j \phi)^2 - m^2 \phi^2 \right). \quad (19)$$

⁸We could weaken this assumption if required. For example, a generalization to Riemannian or Hilbert manifolds of these notions could be obtained by considering that Q is the image of a chart, i.e., that we are working in some local coordinates.

⁹The set of square integrable functions with two (weak) derivatives that are also square integrable.

¹⁰Note how all this includes the description of a field that evolves along Cauchy surfaces Σ_t of a certain foliation of some spacetime $M = \cup_t \Sigma_t$. In particular, since every globally hyperbolic spacetime is globally isometric to an $M := (t_0, t_f) \times \Sigma_0$, we would treat instead of $H^2(\mathbb{R}^3, \mathbb{R})$, a field configuration-space $H^2(\Sigma_0, \mathbb{R})$ so that field paths over spacetime are given by paths on $H^2(\Sigma_0, \mathbb{R})$. In local charts adapted to the foliation, in the image of the charts, we essentially have paths on $H^2(\Omega, \mathbb{R})$ for some $\Omega \subset \mathbb{R}^3$, so the treatment in the text could be seen as a local coordinate picture.

The equations we want to “quantize” are its Euler-Lagrange equations, i.e., conditions defining critical points of the functionals

$$S(q) := \int_{t_0}^{t_f} L(q, \dot{q}) dt \quad \text{and} \quad S(\phi) := \int_{t_0}^{t_f} L(\phi, \dot{\phi}) dt. \quad (20)$$

2. Assuming that the Lagrangian $L(\eta, \dot{\eta})$ has non-degenerate velocity Hessian for given points η ,¹¹ define the canonical conjugate “momentum” map $\chi(\eta, \dot{\eta}) := \text{grad}_{(\cdot, \cdot)}(L(\eta, \cdot))|_{\dot{\eta}}$. In our examples this leaves

$$p_j(q, \dot{q}) = \frac{\partial L(q, \dot{q})}{\partial \dot{q}_j} = m_j \dot{q}_j \quad \text{and} \quad \pi(\phi, \dot{\phi}) = \frac{\delta L}{\delta \dot{\phi}}|_{(\phi, \dot{\phi})} = \frac{\partial \mathcal{L}(\phi, \dot{\phi}, \partial_j \phi)}{\partial \dot{\phi}}|_{(\phi, \dot{\phi}, \frac{\partial \phi}{\partial x_j})} = \dot{\phi}. \quad (21)$$

These make $TQ \simeq Q \times Q$ (since Q is just a vector space) a symplectic vector space with symplectic form $\omega((\eta, \chi), (\eta', \chi')) := \langle \chi, \eta' \rangle - \langle \chi', \eta \rangle$. There, by a Legendre transform we get the Hamiltonian function $H(\eta, \chi) := (\langle \chi, \dot{\eta} \rangle - L(\eta, \dot{\eta}))|_{\dot{\eta}=(\text{grad}(L(\eta, \cdot))|_{\cdot})^{-1}(\chi)}$, e.g.,

$$H(q, p) = \sum_j \frac{1}{2m_j} p_j^2 + V(q) \quad \text{and} \quad H(\phi, \pi) = \int_{\mathbb{R}^3} \frac{1}{2} \left[\pi^2(x) + \sum_j \left(\frac{\partial \phi(x)}{\partial x_j} \right)^2 + m^2 \phi^2(x) \right] d^3x. \quad (22)$$

The solutions to the ODEs we wanted to quantize are then given by the flow map of the symplectic gradient of these Hamiltonians, i.e., their Hamilton’s equations yield the Euler-Lagrange equations.

3. Until here everything can be done rigorously! Now, the (bosonic)¹² quantization recipe tells the following heuristic steps (completely **non-rigorous** for (b) and (c)).

- (a) **If $\eta \in Q$ are functions over a finite set (i.e., a tuple), $\eta = (\eta_1, \dots, \eta_N)$:** Consider for simplicity that $Q = \mathbb{R}^N$, denoting $\eta = q$ and $\chi = p$. The recipe tells to consider the space $L^2(Q, \mathbb{C}, d^N x)$ with $d^N x$ being the Lebesgue measure. There we are told to define some (unbounded) operators \hat{q}_j, \hat{p}_j given for suitable $\psi \in L^2(\mathbb{R}^N, \mathbb{C}, d^N x)$ and $d^N x$ -almost everywhere (a.e.) in q by

$$(\hat{q}_j \psi)(q) = q_j \psi(q) \quad \text{and} \quad (\hat{p}_j \psi)(q) = -i\hbar \frac{\partial}{\partial q_j} \psi(q). \quad (23)$$

It is essential that they satisfy the canonical commutation relations in their joint domain, namely, $[\hat{q}_j, \hat{p}_k] = i\delta_{jk}$, $[\hat{q}_j, \hat{q}_k] = 0$ and $[\hat{p}_j, \hat{p}_k] = 0$. Then, the quantum version of the q system is described by a time evolving wavefunction $t \mapsto \psi_t \in L^2(\mathbb{R}^N, \mathbb{C}, d^N q) \forall t$ satisfying the Schrödinger Equation (SE)¹³

$$(i\hbar \frac{d}{dt} \psi_t)(q) = (H(q, \hat{p}))\psi_t(q). \quad (24)$$

For our driving example’s H we get the familiar SE,

$$i\hbar \frac{d}{dt} \psi_t(q) = - \sum_j \frac{\hbar^2}{2m_j} \frac{\partial^2}{\partial x_j^2} \psi(q) + V(q)\psi(q). \quad (25)$$

Moreover, if $\|\psi_t\|_{L^2}^2 = 1$, the prescription tells us that the probability to find the system in configuration q at time t in a (measurable) set $B \subseteq \mathbb{R}^N$ is given by the probability measure $|\psi(q)|^2 d^N q$, i.e.,

$$\text{Prob}(q \in B | \psi) = \int_B |\psi(q)|^2 d^N q. \quad (26)$$

¹¹If it is degenerate, as turns out to be the case in most cases of interest, then one needs to consider constrained Hamiltonian system techniques. Essentially, one must find a set of constraints in the formally corresponding phase space, that define a submanifold restricted to which the Lagrangian’s gradient does give a bijection. In that submanifold of phase space, after a canonical transformation, we can proceed with the quantization almost as explained here. For several examples and references see (Struyve, 2010). For some constraint types what we get is a so-called “gauge theory”, where the coordinates tangential to the constraint submanifold after the canonical transformation will be directly related to the gauge quantities, elements left underdetermined by the time evolution.

¹²We will mention fermionic particle theories in §B.3.

¹³This is an ill-defined statement: because operator ordering matters, we could define different quantum theories as a result of the quantization by just exchanging variables that before “operatorization” used to commute. Lets ignore this for now.

(b) If $\eta \in Q$ are functions over an uncountably infinite set X , $\eta = (\eta_x)_{x \in X}$: Consider for simplicity $Q = H^2(\mathbb{R}^3, \mathbb{R})$, denoting now $\eta = \phi$ and $\chi = \pi$. We could naively generalize the same prescription and claim that now the relevant space is $L^2(H^2(\mathbb{R}^3, \mathbb{R}), \mathbb{C}, d\mu)$ with $d\mu$ some generalization of the Lebesgue measure $d^N x$ when $N \rightarrow \infty$ (spoiler, there is no obvious choice). Then, we could introduce the operators $\hat{\phi}_x$ and $\hat{\pi}_x$ that for suitable $\Psi \in L^2(H^2(\mathbb{R}^3, \mathbb{R}), \mathbb{C}, d\mu)$, (functions taking value on fields), act μ -a.e. as

$$(\hat{\phi}_x \Psi)(\phi) = \phi(x) \Psi(\phi) \quad \text{and} \quad (\hat{\pi}_x \Psi)(\phi) = -i\hbar \frac{\delta \Psi}{\delta \phi(x)} \Big|_{\phi}. \quad (27)$$

Such operators would formally satisfy the canonical commutation relations “[$\hat{\phi}_x, \hat{\pi}_y$] = $i\delta(x - y)$ ”, [$\hat{\phi}_x, \hat{\phi}_y$] = 0 and [$\hat{\pi}_x, \hat{\pi}_y$] = 0 for $x, y \in \mathbb{R}^3$. Then, according to the recipe, the dynamical equation ruling the evolution of the wavefunctional $t \mapsto \Psi_t(\phi)$ (a path in $L^2(H^2(\mathbb{R}^3, \mathbb{R}), d\mu)$) would be¹⁴

$$i\hbar \frac{d}{dt} \Psi_t(\phi) = H \left(\phi, -i\hbar \frac{\delta}{\delta \phi(x)} \right) \Psi_t(\phi), \quad (28)$$

which leaves for our study case

$$i\hbar \frac{d}{dt} \Psi_t(\phi) = \frac{1}{2} \int_{\mathbb{R}^3} \left[-\hbar^2 \frac{\delta}{\delta \phi(x)} \frac{\delta}{\delta \phi} \Psi|_{\phi} + \left(\sum_{j=1}^3 (\partial_j \phi)^2(x) + m^2 \phi^2(x) \right) \Psi(\phi) \right] d^3 x. \quad (29)$$

Moreover, if $\|\Psi_t\|_{L^2}^2 = 1$, the recipe claims that the system will be found to be a field ϕ in some measurable subset $B \subseteq H^2(\mathbb{R}^3, \mathbb{R})$ following a probability measure $|\psi(\phi)|^2 d\mu$, i.e.,

$$Prob(\phi \in B | \Psi) = \int_B |\Psi(\phi)|^2 d\mu. \quad (30)$$

The generalization to the case in which there are several fields in the Lagrangian, say $\{\phi_1, \dots, \phi_N\}$ would then be straightforward, yielding N field operators $\{\hat{\phi}_1(x), \dots, \hat{\phi}_N(x)\}_{x \in \mathbb{R}^N}$ and $\{\hat{\pi}_1(x), \dots, \hat{\pi}_N(x)\}_{x \in \mathbb{R}^N}$ in an analogous manner to the N particle case.

From the mathematical perspective, there are two main problems with this recipe for generating QFTs. On the one hand, it is not obvious how to define a field of operators like $\hat{\phi}_x, \hat{\pi}_y$ that satisfy the ill-defined condition “[$\hat{\phi}_x, \hat{\pi}_y$] = $i\delta(x - y)$ ”, and take values everywhere in $x \in \mathbb{R}^d$. Most importantly, there exists no obvious generalization of the Lebesgue measure to infinite dimensional spaces. Its characteristic feature was to be the unique measure that is translation invariant in the Borel sigma algebra of \mathbb{R}^n , but no such measure exists if the space is an infinite dimensional Banach space, as we will find later.

We used (a) as a template to come up with (b), but there is a third interesting option we did not consider explicitly yet:

(c) If $\eta \in Q$ are functions over a countably infinite set, $\eta = (\eta_j)_{j \in \mathbb{N}}$: Consider for simplicity $Q = \ell^2(\mathbb{N}, \mathbb{R})$, denoting now $\eta = \alpha = (\alpha_1, \alpha_2, \dots)$ and $\chi = \pi = (\pi_1, \pi_2, \dots)$. Then, we consider a Hamiltonian system of canonical variables (α, π) with infinite dimensional symplectic form as above. Given a Hamiltonian, say,

$$H(\alpha, \pi) = \frac{1}{2} \sum_{j=1}^{\infty} \left(\pi_j^2 + (|\vec{k}_j|^2 + m^2) \alpha_j^2 \right), \quad (31)$$

the generalization of the canonical quantization prescription would tell us that the quantum counterpart of this system lives in the space $L^2(\ell^2(\mathbb{N}), \mathbb{R}, d\mu)$, for some generalization of the Lebesgue measure $d\mu$ (still no obvious choice). Furthermore, there should be some unbounded operators $\hat{\alpha}_j$ and $\hat{\pi}_j$ acting for a suitable $\psi \in L^2(\ell^2(\mathbb{N}), \mathbb{R}, d\mu)$ as multiplication and differentiation operators

$$(\hat{\alpha}_j \psi)(\alpha) = \alpha_j \psi(\alpha) \quad \text{and} \quad (\hat{\pi}_j \psi)(\alpha) = -i\hbar \frac{\partial}{\partial \alpha_j} \psi(\alpha) \quad (32)$$

¹⁴Up to the same problems related to operator ordering as in (a).

and satisfying (in a dense domain) the canonical commutation relations $[\hat{\alpha}_k, \hat{\pi}_j] = i\delta_{kj}$, $[\hat{\alpha}_k, \hat{\alpha}_j] = 0$ and $[\hat{\pi}_k, \hat{\pi}_j] = 0$ (note how there is no longer the same trouble as in the continuum). Then, the time evolution of the quantum system is given by the wavefunctional path $t \mapsto \Psi_t \in L^2(\ell^2(\mathbb{N}), \mathbb{R}, d\mu)$, obeying

$$i\hbar \frac{d}{dt} \Psi_t(\alpha) = H(\alpha, \hat{\pi}) \Psi_t(\alpha), \quad (33)$$

which would read for the example Hamiltonian as

$$i\hbar \frac{d}{dt} \Psi_t(\alpha) \equiv \frac{1}{2} \sum_k \left(-\hbar^2 \frac{\partial^2}{\partial \alpha_k^2} + (m^2 + |\vec{k}_j|^2) \alpha_j^2 \right) \Psi_t(\alpha). \quad (34)$$

Finally, if $\|\Psi\|_{L^2}^2 = 1$, the recipe suggests a probabilistic interpretation by which the system is found in some configuration in a (measurable) set $B \subseteq \ell^2(\mathbb{N})$ with probability measure $|\Psi|^2 d\mu$, namely,

$$Prob(\alpha \in B | \Psi) = \int_B |\Psi|^2(\alpha) d\mu. \quad (35)$$

The problem of having no Lebesgue measure analogue still remains, but in principle the condition on the operators seems a bit more straightforwardly rigorifiable without needing to talk about distributions.

The fun part about (c) is that because we demanded the configuration-spaces for fields over the continuum to be separable Hilbert spaces \mathcal{H} , before doing the quantization we know that there is a unitary 1 to 1 correspondence between such field configuration-spaces \mathcal{H} and ℓ^2 : given a fixed ONB $\{\psi_j\}_{j \in \mathbb{N}} \subset \mathcal{H}$,

$$U : \quad \mathcal{H} \quad \longrightarrow \quad \ell^2(\mathbb{N}) \\ \psi = \sum_{j \in \mathbb{N}} \alpha_j \psi_j \quad \longmapsto \quad (\alpha_1, \alpha_2, \dots). \quad (36)$$

At the pre-quantization step, this allows to map any field configuration-space (of functions over a continuum), into a Hilbert space of functions over countably many elements. This map is arguably given by a canonical transformation, so we get an equivalent pre-quantization Hamiltonian system (the example (31) actually corresponds to the KG Hamiltonian using the unitary for a sine/cosine ONB —we will find it out later). One could view this move as a change of coordinates in infinite dimensions prior to quantization. However, it is not clear at all that the quantization step would then produce equivalent quantum theories using (b) or (c) (if we could make both of them rigorous somehow, as a prerequisite to compare them!).

As a little spoiler, both (b) and (c) will be rigorously doable under suitable cut-offs (finite dimensional approximations), under which one could study this equivalence. However, as the main hit of the present manuscript, we will find that a substantial part of (c) can be built rigorously also in infinite dimensions.

Before all that, let us comment how naturally the heuristic recipes of this section suggest (at least in a heuristic level) a Bohmian-like theory (which we will denote by “pilot-wave theory”).

B.2. How Canonical Quantization Shouts for a Pilot-wave Theory

For many finite dimensional cases (e.g., when the Hamiltonian is a quadratic function of the momentum), one can find that the Schrödinger equation $i\hbar \frac{d}{dt} \psi_t = H \psi_t$ implies the following continuity equation for $|\psi_t|^2(q)$ (obtained by taking $\frac{d\psi_t \psi_t^*}{dt} = \psi_t \frac{d\psi_t^*}{dt} + \psi_t^* \frac{d\psi_t}{dt}$ and plugging the SE there),

$$\frac{d|\psi_t|^2}{dt} = -\vec{\nabla} \cdot (|\psi_t|^2 \vec{v}^{\psi_t}). \quad (37)$$

There, v^{ψ} is given by

$$v_j^{\psi_t}(q, t) := \operatorname{Re} \left\{ \frac{\frac{i}{\hbar} [H(q, \hat{p}), \hat{p}_j] \psi_t(q)}{\psi_t(q)} \right\} \underset{\text{if } H \text{ as (25)}}{=} \frac{1}{m_j} \operatorname{Re} \left\{ \frac{-i\hbar \frac{\partial}{\partial q_j} \psi_t(q)}{\psi_t(q)} \right\} \underset{\substack{\text{wherever polar form} \\ \psi(q) = R_t(q) e^{iS_t(q)/\hbar} \\ \text{is allowed}}}{=} \frac{1}{m_j} \frac{\partial S_t(q)}{\partial q_j}. \quad (38)$$

The interesting part of the continuity equation is that the flow lines of its vector field $v^\psi(q, t)$ are $|\psi|^2 d^n q$ equivariant.¹⁵ It is then natural to postulate the existence of an actual trajectory $Q^\xi(t)$ with “initial” configuration $Q^\xi(t_0) = \xi$ that obeys the flow-line ODE

$$\frac{dQ_k^\xi(t)}{dt} = v^{\psi_t}(Q^\xi(t), t), \quad (39)$$

such that by definition of equivariance, if ξ was “sampled” according to $|\psi_{t_0}(q)|^2 d^n q$, then $Q^\xi(t)$ will be $|\psi_t(q)|^2 d^n q$ distributed at all times t . Such a trajectory (that is always there even if there is no measurement) would explain simultaneously that we always find a fixed configuration (say, point-like particles) in a measurement (never a superposition) and why still we find a stochastic outcome following the Born rule. For this, one postulates the “quantum equilibrium hypothesis”, as brilliantly explained by Dürr et al. (1992).

By the flow equation above, one then says that $Q^\xi(t)$, the actual configuration of the system in time, is “piloted” by the configuration-space wave $\psi_t(q)$ and we talk about a *pilot-wave theory*. Note that in modern treatments, the “piloting wave” is not considered to be something that materially “pushes” the ontological configuration, but rather as part of a dynamical law of motion. That is, it is seen as a *nomological* element (nobody asks “where” the piloting wave is or whether it is observable, in the same way as nobody used to ask “where” Newton’s second laws were or whether they were observable).¹⁶

With the very same spirit, **formally**, one can find a continuity equation underlying many relevant infinite dimensional configuration-space Schrödinger equations, as exemplified by Struyve (2010). In particular, one can find a heuristic continuity equation for our KG study case. In direct analogy with the finite dimensional case, one could then take the continuity equation and obtain an equivariant vector field $v^{\Psi_t}(\phi, t)$, allowing the definition of some equivariant flow-lines over field configuration-space (where each trajectory is now a time evolving **field**). The pilot-wave picture would then be completed by declaring that there exists an actual field of trajectory $t \mapsto \Phi^\Xi(t)$, which is the flow-line going through some initial configuration $\Phi^\Xi(t_0) = \Xi$. Namely, the solution of

$$\frac{d\Phi^\Xi(t)}{dt} = v(\Phi^\Xi(t), t); \quad \Phi^\Xi(t_0) = \Xi. \quad (40)$$

Then, if Ξ was sampled according to $|\Psi_{t_0}(\phi)|^2 d\mu$, $\Phi^\Xi(t)$ will be $|\Psi_t(\phi)|^2 d\mu$ distributed at all times t and one could perhaps propose a similar analysis as Dürr et al. (1992) to give a consistent interpretation to the theory. We would say then that “the actual field is piloted” by the wavefunctional and call it a *pilot-wave theory*. The case of N fields in the Lagrangian would likewise yield N ontological physical space fields, all piloted by the same wavefunctional.

With all, we could give a (rather informal) definition of what we mean by a pilot wave theory:

Definition 1. A **pilot-wave theory** (PWT) is the tuple formed by:

- a measure space (Q, μ) that is (ideally) a normed vector space (NVS) (and more ideally an inner product vector space (IPVS)), called *configuration-space*,
- a Hilbert space $\mathcal{H} := L^2(Q, \mathbb{C}, d\mu)$
- a (densely defined) self-adjoint operator on \mathcal{H} called Hamiltonian \hat{H} generating the dynamics for the vectors $\Psi_0 \in \mathcal{H}$, called wavefunction(al)s, via a Schrödinger equation: $i \frac{d}{dt} \Psi_t = \hat{H} \Psi_t$ (or, equivalently by Stone’s theorem, a SCOPUG¹⁷ $\{U_t = e^{-i \frac{t}{\hbar} \hat{H}}\}_t$)
- for each path in \mathcal{H} solving the SE $t \mapsto \Psi_t \in \mathcal{H}$, a vector field v^{Ψ_t} on $Q \times \mathbb{R}$, whose flow-map is $|\Psi_t| d\mu$ -equivariant
- the assumption that for each “physical realization of Ψ_t ”, the primitive ontology is given by a single trajectory on Q . It is epistemologically unknown but ontologically determined. It is one of the flow-lines

¹⁵Essentially meaning that if $(\xi, t) \mapsto Q(\xi, t)$ is the flow map of the velocity field v^{ψ_t} , i.e., $\frac{\partial}{\partial t} Q(\xi, t) = v^\psi(Q(\xi, t), t)$ (the trajectories are tangent to the velocity field) with initial condition $Q(\xi, t_0) = \xi$, then for all time t and (measurable) $B \subseteq \mathbb{R}^n$,

$$\int_B |\psi_{t_0}|^2(q) d^n q = \int_{Q(B, t)} |\psi_t|^2(q) d^n q.$$

¹⁶Still, there are ways using the concept of conditional wavefunctions and infinite chains of potential energy fields, as explained by Norsen et al. (2015), to build a physical-space field ontology for the wavefunction within the scope of the standard SE (25).

¹⁷Strongly continuous one parameter group of unitary operators.

of v^Ψ , with an epistemological measure of typicality such that it crosses a subset $B \subseteq Q$ with a probability $\int_B |\Psi_t|^2 d\mu$. ♦

(Note how a priori such a pilot wave theory does **not** bring with it a connection between the ontological trajectory on Q and our phenomenological world —nor what a “physical realization of Ψ_t ” means. Thus, it is not yet a complete statement of a physical theory!).

B.2.2. Digression: Why is the quantization of fields called “second quantization”?

“Aren’t we doing only one quantization?” Well, we could do the same procedure as for the KG equation and charge density, now for the Schrödinger equation and a wavefunction.¹⁸ Then, the result would be a wavefunctional $\Psi(\psi)$ over wavefunctions ψ . That is, a wavefunctional, which is already the quantization of a set of particles, would be “quantized again” to get a wavefunctional of wavefunctions.

Certainly, if we take seriously the term “second quantization”, one could (misguidedly) end up thinking that $|\Psi(\psi)|^2 d\mu$ is the probability measure to find a wavefunction ψ . And that “the outcome ψ ”, via $|\psi(q)|^2 d^N q$, yields the probability measure to find the particles in configuration q . Even worse, the obvious field ontology would tell us that there is an actual wavefunction ψ at all times, which could then have a subordinate particle ontology, giving a Bohmian trajectory at all times. Moreover, the ontological wavefunction would happen to follow the Schrödinger equation except for an additional inhomogeneity due to the “quantum potential” (analogous to the inhomogeneity we got for the Maxwell equation in the introduction). Hence, it would no longer possess exactly a unitary time evolution. And then, a myriad of speculative (and possibly misguided) questions would emerge: what if the extra term is exactly like the inhomogeneity we get in the dynamical equation of a conditional wavefunction (a slice of a wavefunction)? Would it mean there are additional spatial degrees of freedom and our Universal wavefunction is a slice of it?

To make it even better, once such a “second quantization” is considered, we could speculate about a “**third** quantization”. We could now take a wavefunctional-**al** $\mathfrak{P}(\Psi)$ over the space of possible wavefunctionals Ψ (which live in turn over the space of wavefunctions ψ , which live in turn over some particle configuration space Q). Then we would postulate some “field-al” operators that act as multiplication and functional derivative with respect to wavefunctionals, give a SCOPUG/“meta-”Schrödinger equation and again we could give it a field ontology. But then, why stop here? We could consider a fourth quantization and a fifth quantization and a sixth one, until we get sufficiently mad! Would there be any limit in reasonable terms for such an iteration? What kind of new experimental predictions would it give?

Please, do not take this parenthetic section seriously, its objective was merely to make some comedy while explaining the misleading origin of the name.

B.3. What about Fermions?

The quintessential relativistic field describing a fermion is the Dirac 4-spinor wavefunction $\psi \in L^2(\mathbb{R}^3, \mathbb{C}^4)$. Although it has its origin in quantum mechanics, it can also be understood as a continuous charge distribution within classical mechanics, with several polarization coordinates or fields required to specify the charge field’s dynamics (hence the several components in each spatial position). Then, we could proceed with the second quantization also for this field and find a wavefunctional of Dirac 4-spinors $\Psi(\psi)$ making a second quantized QFT out of it. However, there is a key difficulty that makes the field configuration-space approach very unappealing for fermions.

As there are four spinor components, namely, four fields that make up what we call the Dirac field (if we would like to make them complex then eight), we would need to consider (at least) four field operators $\{\hat{\psi}_1(\vec{x}), \hat{\psi}_2(\vec{x}), \hat{\psi}_3(\vec{x}), \hat{\psi}_4(\vec{x})\}_{\vec{x} \in \mathbb{R}^3}$. As before, the obvious generalization of canonical quantization would demand the field operators to act as multiplication operators on the wavefunctional:

$$(\hat{\psi}_j(\vec{x})\Psi)[\psi, t] = \psi_j(\vec{x})\Psi(\psi) \quad \forall j \in \{1, 2, 3, 4\}; \forall \vec{x} \in \mathbb{R}^3, \quad (41)$$

¹⁸By perhaps taking the real and imaginary parts as independent fields, since after all we are considering Hamiltonian systems over *real* vector spaces

together with an anticommutation relation, including among others,¹⁹

$$\{\hat{\psi}_i(\vec{x}), \hat{\psi}_j(\vec{x})\} = 0 \quad \forall i \neq j \in \{1, 2, 3, 4\}, \forall \vec{x} \in \mathbb{R}^3. \quad (42)$$

The two above conditions together mean that for relevant $\Psi \in L^2(L^2(\mathbb{R}^3, \mathbb{C}^4), \mathbb{C})$,

$$0 = \{\hat{\psi}_i(\vec{x}), \hat{\psi}_j(\vec{x})\} \Psi(\psi) = (\psi_i(\vec{x})\psi_j(\vec{x}) + \psi_j(\vec{x})\psi_i(\vec{x}))\Psi(\psi). \quad (43)$$

But because complex numbers commute, this is only possible if either $\psi_j(\vec{x})$ or $\psi_i(\vec{x})$ are zero. Hence, we are forced to consider a space of spinor wavefunctions $\psi \in L^2(\mathbb{R}^3, V^4)$ that do not take values in \mathbb{C} , but some other “number” space V where numbers anticommute. This is mathematically realizable as so-called *Grassman numbers*.

However, there are several issues with Grassman valued Dirac fields, starting from difficulties to define for each field in such a configuration space concepts like a charge or energy density, which will not be real-valued and hence are problematic for pre-quantization electromagnetic interactions. But most relevantly after quantization, there are (even more) difficulties to interpret the wavefunctional’s amplitude-squared as a probability density and to give a measure to the space of Grassman valued fields. Moreover, it is a common practice to use Grassman valued wavefunctionals instead of complex valued ones, which makes the problem even worse: probability amplitudes would not even be real valued... (See [Struyve \(2010\)](#); [Sebens \(2022\)](#) for more details of the topic and how to still approach fermionic QFTs within a field ontology picture in a heuristic way.)

As opposed to all these problems, if we treat fermions instead with a particle ontology in mind, very appealing QFTs are obtained: ranging from Fock spaces, till infinite wedge product spaces or time varying Dirac seas ([Dimock, 2011](#); [Deckert et al., 2010](#)). Moreover, the cut-off picture there (in particular, assuming space is finite and a UV cut-off exists) allows to employ a theory with a fixed number of fermions to describe all QFTs, as exposed by [Deckert et al. \(2016\)](#). That is, no creation and annihilation of fermions would need to happen ontologically, and they would only be an effective way to talk about fermions relative to the “sea level” of such a finite Dirac sea. With all, we will not treat fermions any more in the context of field ontologies.

It is remarkable that diametrically opposed to the complication of describing a field ontology for fermions, there are problems to give a particle ontology to bosons. Specifically, to those bosons that lack a clear “single particle” quantum theory, such as photons —a lack that is explained for instance by [Tumulka \(2022\)](#).

All in all, the moral is that a mixed ontology appears to be the nicest one in mathematical terms: describe fermions as particles (ontologically) while bosons as physical-space fields. They are “virtual particles”, in the sense that their particle aspects are rather due to “bumps” of some ontological field spreading in physical space. This turns out to match with the way we trained ourselves to picture the world in classical mechanics, so that is a good news Or...should we be worried instead? The preference of the maths for such a mixed theory could be interpreted instead as our own preference (or its historical drift), propagated to the developed mathematical formulations. That is, as a symptom of the biased way in which we built the foundations of our physical theories, which perhaps unwillingly shaped a mathematical trap with no easy escape.

¹⁹Instead of commutation, due to the different statistics for identical fermions —essentially Pauli’s exclusion principle.

C. RIGOROUS FIELD-CONFIGURATION QM I: MODE EXPANSION CUT-OFF

In sections C and D, we present the two main paths towards solving the mathematical ill-definitions we spotted in the canonical quantization of fields. The first way, explained in C, is to impose enough conditions to the configuration space that only elements with finitely many degrees of freedom within it become relevant. Then we recover the finite dimensional scenario and both raised problems instantly vanish: the one about the operators and the one about the measure.

Since we assumed the field configuration space to be a separable Hilbert space, there exists a countable orthonormal basis (ONB) so that elements of such a configuration-space are in unitary one-to-one relation with elements of $\ell^2(\mathbb{N})$. Thus, before the quantization we could express the Hamiltonian system over the space of possible expansion coefficients of some ONB, $\ell^2(\mathbb{N})$, (likewise for their canonically conjugate momenta). A quantization in this (apparently) equivalent space would likely erase the problem of field operators $\hat{\phi}_x$ not being easily definable for all $x \in \mathbb{R}^3$ (now the indexing set for field operators would be countable, see §B.1.c). Nevertheless, we would still have the problem of an infinite dimensional vector space, with no distinguished choice of “uniform” measure to build the wavefunctional space of.

To solve the measure’s issue, before quantization we can make an assumption. We assume that the only fields in the configuration space that matter are the ones having an expansion where the coefficients that come after the M -th ONB element are zero. This then leaves a finite dimensional configuration space (namely, the truncated expansion coefficient space) where the Lebesgue measure is available again.

Hence, we need two ingredients to rigorify the second quantization of some field theory: (a) an ONB of the field configuration-space and (b) a point of truncation. Both choices should be guided by physical intuition.

One of the most commonly employed such physical assumptions is the following one. Arguably, “the spatial variations of a physical-space field below some very insignificant characteristic length-scale $l > 0$ are meaningless”. This can be interpreted in two different and closely related ways (yielding different ONB and truncation points as candidates).

C.1. Ultraviolet (UV) Cut-off

One way to understand the limitation on the spatial variation is as a limitation of the “frequency” that sinusoids composing the field could have in a Fourier basis. (This is precisely the trick we used to make the introductory example sound.)

Since plane-waves are not elements of a space like $H^2(\mathbb{R}^3, \mathbb{R})$, to use this trick we first need another assumption: that space is finite,²⁰ say $\Omega := (0, L)^3$, so that plane-waves are square integrable and there is for instance an ONB of cosines $\left\{ \sqrt{\frac{2}{L^3}} \cos(\langle \vec{k}_j, \cdot \rangle_2) \right\}_{j \in \mathbb{N}}$. Given that, for any $\phi \in H^2(\Omega, \mathbb{R})$ there exist coefficients $\alpha_j \in \mathbb{R}$ such that

$$\phi(\vec{x}) \equiv \sum_{j \in \mathbb{N}} \alpha_j \sqrt{\frac{2}{L^3}} \cos(\vec{k}_j \cdot \vec{x}). \quad (44)$$

(Also re-writable as a complex exponential series with real coefficients²¹). Now yes, we can implement that only spatial variations above the length-scale l are significant by taking only \vec{k}_j for which the wavelength of the cosine is bigger than l , i.e., $2\pi/|\vec{k}_j| > l$. That is, we impose $|\vec{k}_j| < 2\pi/l$, cutting-off the plane-waves at a finite, M -th frequency.

Defining $u_j := \sqrt{\frac{2}{L^3}} \cos(\langle \vec{k}_j, \cdot \rangle)$, the truncated expansion of (44) reads, for a trajectory $t \mapsto \phi_t \in H^2(\Omega, \mathbb{R})$, as $\phi_t = \sum_{j=1}^M \alpha_j(t) u_j$ with $\dot{\phi}_t = \sum_{j=1}^M \dot{\alpha}_j(t) u_j$. If we plug them in our driving case’s pre-quantization

²⁰Justifiable because we might be considering a cavity as before, or because the space-slice we are considering is (perhaps big but) finite in extent, as is the case in several relevant cosmological spacetimes models for the Big-Bang.

²¹Defining $\alpha_{-j} := \alpha_j$ and $c_j := \sqrt{\frac{1}{2L^3}} \alpha_j$, then:

$$\phi(\vec{x}) \equiv \sum_{j \in \mathbb{N}} \alpha_j \sqrt{\frac{2}{L^3}} \cos(\vec{k}_j \cdot \vec{x}) = \sum_{j \in \mathbb{N}} \alpha_j \sqrt{\frac{1}{2L^3}} (e^{i\vec{k}_j \cdot \vec{x}} + e^{-i\vec{k}_j \cdot \vec{x}}) = \sum_{j \in \mathbb{Z}} c_j e^{i\vec{k}_j \cdot \vec{x}}$$

Hamiltonian (22), and we apply the orthonormality of the basis we get

$$H(\phi(\alpha_j), \pi = \dot{\phi}(\alpha_j)) = \frac{1}{2} \sum_{j=1}^M \left(\dot{\alpha}_j^2 + (|\vec{k}_j|^2 + m^2) \alpha_j^2 \right). \quad (45)$$

One can find that this was essentially a canonical transformation if $\dot{\alpha}_j$ is taken to be the canonical momentum of α_j , which we denote by π_j . Hence, at the classical level, the Hamiltonian system is still the same (albeit in different coordinates). Presumably the resulting quantized theory will also need to be the same, but we will not care about this here.

Now, note that the Hamiltonian we obtained,

$$H(\alpha_j, \pi_j) = \frac{1}{2} \sum_{j=1}^M \left(\pi_j^2 + (|\vec{k}_j|^2 + m^2) \alpha_j^2 \right) \quad (46)$$

matches exactly with the Hamiltonian of M uncoupled harmonic oscillators with positive spring coefficients depending on the frequency of the represented frequency-mode. Using the unproblematic quantization of finitely many degrees of freedom, we get that the quantized system is described by a wavefunction $\Psi \in L^2(\mathbb{R}^M, \mathbb{R}, d^M x)$, where the Lebesgue measure can be used with no problem, solving the main problem we had. On the other hand, we get the well-defined ‘‘field operators’’ with dense domain (e.g., the Schwartz functions), defined as $\hat{\alpha}_j \Psi_t(\alpha_1, \dots, \alpha_M) := \alpha_j \Psi(\alpha_1, \dots, \alpha_M)$ and $\hat{\pi}_j \Psi_t(\alpha_1, \dots, \alpha_M) := i\hbar \frac{\partial}{\partial \alpha_j} \Psi(\alpha_1, \dots, \alpha_M)$. Denoting $\alpha := (\alpha_1, \dots, \alpha_M)$, the wavefunction obeys the Schrödinger equation

$$i\hbar \frac{d}{dt} \Psi_t(\alpha) \equiv \sum_k \left(-\frac{\hbar^2}{2} \frac{\partial^2}{\partial \alpha_j^2} + (m^2 + |\vec{k}_j|^2) \right) \Psi_t(\alpha). \quad (47)$$

By the Born rule, $|\Psi(\alpha)|^2 d^N \alpha_j$ gives the probability measure for finding the field in a configuration $\sum_{j=1}^M \alpha_j u_j$. We find the continuity equation for its density $|\Psi(\alpha)|^2$, to be given by

$$\frac{d}{dt} |\Psi_t|^2(\alpha) = -\vec{\nabla} \cdot (|\Psi(\alpha)|^2 \vec{v}^\Psi(\alpha)) \quad (48)$$

where,

$$\vec{v}_t^\Psi(\alpha) = \frac{1}{m} \operatorname{Re} \left\{ \frac{-i\hbar \vec{\nabla} \Psi_t(\alpha)}{\Psi(\alpha)} \right\} \stackrel{\text{wherever polar form allowed}}{\Psi_t(\alpha) = R_t(\alpha) e^{iS_t(\alpha)/\hbar}} \stackrel{=}{=} \frac{1}{m} \vec{\nabla} S_t(\alpha). \quad (49)$$

Then we can introduce an ontological field’s trajectory, parametrized by a time evolving tuple of expansion coefficients $\alpha^\xi(t)$ with initial condition $\alpha^\xi(t_0) = \xi$, as the flow-line of the probability-equivariant velocity field:

$$\frac{d}{dt} \alpha^\xi(t) = \frac{1}{m} \vec{\nabla} S(\alpha^\xi(t)). \quad (50)$$

Finally, it might be interesting to mention for our coming discussions, that the ground state for the given Hamiltonian (46) is a Gaussian (since each degree of freedom is just a harmonic oscillator):

$$\Psi(\alpha_1, \dots, \alpha_M) = \mathcal{N} e^{-\sum_{j=1}^M \frac{\alpha_j^2}{2} \frac{\sqrt{m^2 + |\vec{k}_j|^2}}{\hbar}} \quad (51)$$

with some normalization constant \mathcal{N} . And again, using the Hermite polynomials and the Fock representation, one could talk about ‘‘Klein-Gordon-on’’ or ‘‘charge-on’’²² particles of frequency \vec{k}_j , as bosonic (virtual) particles whose number is given by the excitation level of the j -th harmonic oscillator. From there, one could build a particle ontology for the KG equation’s quantization using the Fock space representation.

²²If we are describing a charged continuum.

C.1.1. Digression: The Quantum Potential as the Origin of Mass

Using the polar form of Ψ_t , one can find that the Schrödinger equation implies two real PDEs. If we write one of them from the perspective of a trajectory solving (50) (i.e., in the ‘‘Lagrangian frame’’), we get that the coefficients obey the harmonic oscillator’s Newton’s laws, but with an additional potential term Q^Ψ (the so-called ‘‘quantum potential’’):

$$\frac{d^2}{dt^2}\alpha_j^\xi(t) = \sqrt{m^2 + |\vec{k}_j|^2} \alpha_j^\xi(t) - \frac{\partial}{\partial \alpha_j} Q^\Psi(\alpha, t)|_{\alpha^\xi(t)}. \quad (52)$$

Putting $\Phi(t) = \sum_{j \in \mathbb{N}} \alpha_j(t) u_j$ (by adding the previous equation for $j \in \{1, \dots, M\}$), one can find that this equation implies the ontological field starting at $\Phi^\xi(t_0) = \sum_j \xi_j u_j$ obeys

$$\ddot{\Phi}_t^\xi = \Delta \Phi_t^\xi - m^2 \Phi_t^\xi - F(Q^\Psi)|_{\Phi_t^\xi} \quad (53)$$

for some fixed function F of the quantum potential. That is, it obeys the Klein-Gordon equation but with an extra term giving the deviation of our ontological field’s evolution from a ‘‘non-quantized’’ KG field.

As a funny observation by [Holland \(1995\)](#), if we took a massless particle’s KG equation (i.e., set $m = 0$), the result would have been that the quantized version of it, in the pilot wave theory, would follow an equation

$$\ddot{\Phi}_t^\xi = \Delta \Phi_t^\xi - F(Q^\Psi)|_{\Phi_t^\xi}. \quad (54)$$

If by accident (or not) $F(Q^\Psi)|_{\Phi^\xi} \simeq C \Phi^\xi$ for some constant $C > 0$, then the quantum version of the massless KG equation would be a massive KG equation. The speculative part: could it be that the mass of massless particles is originated in the quantum potential?

C.2. Discrete Space

The other way to understand the limitation on the spatial variation of fields below length-scale l is to assume there is actually no field in between spatial points separated by that distance l . Perhaps because in between those points there is no space at all in the ontological sense, or perhaps because an average/coarse graining per each such tiny neighbourhood gives enough information about the field. That is, we assume that fields take value only over a discretized space of points $\mathcal{L} := \{\vec{x}_j\}_{j \in \mathbb{N}} \subset \mathbb{R}^3$ with grid-spacing l . The discretization of space promotes the formal Dirac deltas into a rigorous ONB (analogous to how making the space finite made the plane waves a rigorous ONB in §C.1). This solves the problem of field operators in a very natural way, because unlike in the UV-cut-off case, we can now define them in position (and frequency) space as $\{\hat{\phi}_x, \hat{\pi}_x\}_{x \in \mathcal{L}}$ (closer to the heuristic $\{\hat{\phi}_x, \hat{\pi}_x\}_{x \in \mathbb{R}^3}$).

So we now have the ingredient (a). But ‘‘ $H^2(\{\vec{x}_j\}_{j \in \mathbb{N}}, \mathbb{R}, d\nu)$ ’’ is still an infinite dimensional vector space (each field still has countably infinitely many degrees of freedom), so we need a cut-off.

Hence, we cut-off the Dirac delta basis at some M -th grid-point, which is equivalent to assume that space is finite! Since this leaves a finite dimensional configuration space, we can then rigorously have wavefunctionals over fields $\Psi(\phi) := \Psi(\phi_{\vec{x}_1}, \dots, \phi_{\vec{x}_M})$, where $\Psi \in L^2(\mathbb{R}^M, \mathbb{C}, d^M x)$ is equipped with the Lebesgue measure. Finally, we can define field operators as we wanted,

$$(\hat{\phi}_{\vec{x}_j} \Psi)(\phi_{\vec{x}_1}, \dots, \phi_{\vec{x}_M}) = \phi_{\vec{x}_j} \Psi(\phi_{\vec{x}_1}, \dots, \phi_{\vec{x}_M}) \quad \text{and} \quad (\hat{\pi}_{\vec{x}_j} \Psi)(\phi_{\vec{x}_1}, \dots, \phi_{\vec{x}_M}) = -i\hbar \frac{\partial}{\partial \phi_{\vec{x}_j}} \Psi(\phi_{\vec{x}_1}, \dots, \phi_{\vec{x}_M}). \quad (55)$$

The Schrödinger equation ruling the motion of $t \mapsto \Psi_t \in L^2(\mathbb{R}^M, \mathbb{C}, d^M x)$ in our study case would be the discretization of the heuristic formula (29),

$$i\hbar \frac{d}{dt} \Psi_t(\phi) = \frac{1}{2} \sum_{j=1}^M \left[-\hbar^2 \frac{\partial^2}{\partial \phi_{\vec{x}_j}^2} \Psi_t(\phi) + \left(\sum_{k=1}^3 \left(\frac{\phi_{\vec{x}_j + \vec{e}_k} - \phi_{\vec{x}_j}}{l} \right)^2 + m^2 \phi_{\vec{x}_k}^2 \right) \Psi_t(\phi) \right], \quad (56)$$

with $\vec{e}_k \in \mathbb{R}^3$ the canonical basis vector in direction k .

Lastly, the probability to find the system in a field $\phi := (\phi_{\vec{x}_1}, \dots, \phi_{\vec{x}_M})$ at time t would be given by the probability measure $|\Psi_t(\phi)|^2 d^M x$. Again, one can then find a continuity equation for $|\Psi_t(\phi)|^2$ in the usual

way, and introduce the trajectory of an actual ontological field $\Phi^\Xi(t)$ starting in $\Phi^\Xi(t_0) = \Xi$. Through the obvious equivariant velocity field behind the continuity equation, one could obtain an equation of motion for each component $\Phi_{x_j}^\Xi$ with $j \in \{1, \dots, M\}$, as a system of ODEs.

This is the rigorization that is closest to what physicists mean by their formal manipulations. In fact, since any PDE in a computer will need to be discretized to a finite lattice, one could naively believe that in the end, either using a discretized version of the theory or a rigorous infinite dimensional rigorification (as we are going to approach in D), the result will not be distinguishable while we use computers to obtain verifiable predictions. This would make attempts to rigorify infinite dimensional formulas rather unmotivated. However, although it might be the case that the infinite dimensional case should coincide with this cut-off approach, if we really had a rigorous theory in infinite dimensions, employing analytic manipulations, predictions of very general kinds could be deduced. Predictions that may even be unreachable through computational manipulations of the discretization (as often happens in mathematical physics).

Finally, by the discrete Fourier transform (DFT) any discretized function on finitely many points, say, the fields over the lattice above, can be written down as finite series over cosines. Then at first glance, the two explicit prescriptions we gave to understand the length-scale cut-off seem to be essentially the same. Both end up requiring space to be finite and end up truncating a sinusoid expansion. It is an open question if they really yield the same theory for common choices of l .

C.3. QFT and the Endless Tower of Approximate Theories

Physicists routinely employ heuristic and ill-defined formulas like

$$\int_{H^2(\mathbb{R}^3, \mathbb{R})} |\Psi(\phi)|^2 \mathcal{D}\phi = \int_{H^2(\mathbb{R}^3, \mathbb{R})} |\Psi(\phi)|^2 \prod_{x \in \mathbb{R}^3} d\phi_x \quad (57)$$

where $\mathcal{D}\phi$ or $\prod_{x \in \mathbb{R}^3} d\phi_x$ pretend to represent a generalization of the Lebesgue measure to infinite dimensional spaces, something like $\lim_{N \rightarrow +\infty} \prod_{j=1}^N dx_j$. In the rigorous discretization we presented, such an expression would essentially be

$$\int_{\mathbb{R}^M} |\Psi(\phi)|^2 \prod_{x_j \in \{x_k\}_{k=1}^M} d\phi_{x_j}, \quad (58)$$

which, unlike (57), is a well-defined expression. It is likely that so many formal manipulations in the physics literature give somewhat meaningful results even if they employ ill-defined equations like (57), precisely because one can understand them as symbols that are meant to represent a discrete expression like (58).

Some branches of QFT have directly abandoned even the formal manipulations in infinite dimensions and have started to assume the discretization of space seriously, as lattice gauge theory exemplifies.

In fact, one should notice that as of today, the only way to make the Standard Model a serious/mathematically meaningful theory is by the cut-off technique. Despite that, remarkably, very precise experimental predictions have been made and confirmed (or perhaps “fitted”...).

The problem with cut-off theories is that, sure, we can make the cut-off as insignificant as we wish...but each choice yields a different theory than the previous one! [Wallace \(2006\)](#) and references therein argue that this *per se*, should not be regarded as a problem. Similar techniques are employed in solid state physics to make sense of functional space integrals and there, where experimental results to contrast calculations are remarkably more accessible than in particle physics, it is found that as long as we choose the cut-off small enough relative to the scale of distances we are interested in, the coarse grained predictions will be independent of the particular cut-off (say, the particular grid employed). Since in solid state physics we know that assumptions about the continuum are really an approximation of a deeper theory in which there are finitely many discrete atoms and molecules, spatial cut-offs are moreover positively motivated.

As [Wallace \(2006\)](#) poses very nicely, the analogue in QFT is somewhat more convoluted. Since the discretization in this case is not an approximation of an underlying truly discrete setting, but of an ideally continuous one, there seems to be no lower bound for the step-size of the discretization. When we are interested in shorter length-scale predictions, we should do the discretization finer and finer. However, since there exists no “underlying” theory that works (yet) at some deeper scale, this makes QFT an endless tower of approximate theories,

each less approximate than the previous one, but nevertheless only valid in a certain coarse grained scale. But then, QFT's current status is that of an intrinsically approximate theory! What a thing! This either tells us to redefine the concept of physical theory, or to wait for a deeper theory to be found, such as a rigorous infinite dimensional version of QFT, or, if that happens to be impossible (as all efforts have shown so-far for interacting theories in 4 dimensional spacetime), perhaps a radically new theory (say, on the line of string theory). In this situation, one is left to think until which point is not QFT merely a fitting technique, as those employed in machine learning. Put enough weights to be fixed by experimental observations and the theory will be able to predict plenty of other experimental observations. We will come back to this in the concluding remarks.

D. RIGOROUS FIELD-CONFIGURATION QM II: INSIST ON INFINITE DEGREES OF FREEDOM

Let us break the pessimistic picture in which a field ontology can only be rigorously sustained under cut-off conditions. Let us look towards a more serious rigorification of the heuristic QFT of section B.1, where by “serious” we mean that we insist on the **infinite** dimensional configuration space. In facing this we find two main issues to address.

D.1. Problem I: The Operator Fields

If we want to take seriously the commutation relation we wanted to impose for the field operators, “ $[\hat{\phi}_x, \hat{\pi}_y] = i\delta(x - y)$ ”, we already see that we require the field operators to be something more related to distributions than proper operators indexed by an infinite set. Now, although there is strictly speaking no absolute impossibility proof of having a well-defined QFT with field operators like $\{\hat{\phi}_x\}_{x \in \mathbb{R}^3}$, there are several no-go theorems under quite general assumptions that discard this option. Halvorson and Müger (2006) in section 6 nicely review some relevant ones. As an example, here we give one:

Proposition 1. Let $\hat{\phi}(x)$ and $\hat{\pi}(y)$ be self-adjoint field operators, either bounded or unbounded, such that $[\hat{\phi}(x), \hat{\pi}(y)] = 0$ when $x \neq y$.²³ If $y \mapsto U(y)$ is a continuous unitary representation of the translation group such that $U(y)\hat{\pi}(x)U(y)^* = \hat{\pi}(x + y)$ for all $x, y \in \mathbb{R}^3$, then $[\hat{\phi}(x), \hat{\pi}(x)] = 0$ for all $x \in \mathbb{R}$. \blacklozenge

But this is in flagrant contradiction with our aimed commutation relation (as we wanted $[\hat{\phi}(x), \hat{\pi}(y)]$ to “blow up” when $x \rightarrow y$). Certainly, we would like to have a unitary that tells us how to change a wave-vector when we change our position-space coordinate system by a translation (hence, a representation of the translation group). But if we want the field operators to change accordingly, we cannot have $\{\hat{\phi}_x\}_{x \in \mathbb{R}^3}$.

Of course, as Halvorson and Müger (2006) mention, we could always weaken some of the hypothesis of a no-go theorem, say, the continuity of the representation in the above case. Then, finding example theories, we could get “go theorems”. After all, history has shown in remarkable occasions, that “a no-go theorem is rather lack of imagination”,²⁴ as Bell pointed out. The general trend in mathematical approaches to QFT is different however. They tell to consider **operator valued distributions**. Instead of $\{\hat{\phi}_x\}_{x \in \mathbb{R}^3}$,²⁵ we define, for (a complexification of) the field configuration-space \mathcal{H} , a set $\{\hat{\phi}(f)\}_{f \in \mathcal{H}}$, with $\hat{\phi}(f)$ heuristically corresponding to

$$\hat{\phi}(f) = “ \int_{\mathbb{R}^3} f(x) \hat{\phi}_x dx ”, \quad (59)$$

i.e., a “smearing” of the field operator by an envelope field f . Intuitively, if we take f to be compactly supported in a small enough ball around some $\vec{x} \in \mathbb{R}^3$, we will properly approximate the idea behind the formal symbol $\hat{\phi}_{\vec{x}}$, $\vec{x} \in \mathbb{R}^3$. For operator valued distributions, on the other hand, the commutation relation would look like

$$[\hat{\phi}(f), \hat{\pi}(g)] = i\langle f, g \rangle Id, \quad (60)$$

with Id the identity on the space where the field operators act. This is a more tractable object in the construction of a mathematically well-posed theory.

²³In the unbounded case, it is meant that $\hat{\phi}(x)$ and $\hat{\pi}(y)$ are defined on a common dense subset D of the Hilbert space where the operators act, and that the field operators commute there.

²⁴No-go theorems have often revealed more about the limitations of current frameworks than about absolute impossibilities...

²⁵Apparently, the only known way to make precise these operator fields over $x \in \mathbb{R}^3$ is by considering them to be sesquilinear form valued maps, with the problem that then there are problems to know what precisely is meant by powers of the field operators. Yet, in the (heuristic) presentations of interacting theories, powers of field operators are the daily grind.

In the particular case of field operators acting on a wavefunctional space like $L^2(\mathcal{H}, \mathbb{C}, d\mu)$ ($d\mu$ yet to be specified), consider $\Psi \in L^2(\mathcal{H}, \mathbb{C}, d\mu)$. Instead of the problematic

$$(\hat{\phi}_x \Psi)(\phi) = \phi(x)\Psi(\phi) \quad \text{and} \quad (\hat{\pi}_x \Psi)(\phi) = -i\hbar \frac{\delta \Psi}{\delta \phi(x)} \Big|_{\phi}, \quad (61)$$

we will have

$$(\hat{\phi}(f) \Psi)(\phi) = \langle f, \phi \rangle \Psi(\phi) \quad \text{and} \quad (\hat{\pi}(f) \Psi)(\phi) = -i\hbar \langle f, \frac{\delta \Psi}{\delta \phi} \Big|_{\phi} \rangle, \quad (62)$$

which in our driving example, $\mathcal{H} = H^2(\mathbb{R}^3, \mathbb{R})$, would just be

$$(\hat{\phi}(f) \Psi)(\phi) = \left(\int_{\mathbb{R}^3} f(x) \phi(x) d^3x \right) \cdot \Psi(\phi) \quad \text{and} \quad (\hat{\pi}(f) \Psi)(\phi) = -i\hbar \int f(x) \frac{\delta \Psi}{\delta \phi(x)} \Big|_{\phi} d^3x. \quad (63)$$

Alternatively, note that if the infinite dimensional configuration-space were, say, $\ell^2(\mathbb{N}, \mathbb{R}, \nu)$, the heuristic expression for field operators, $[\hat{\phi}_k, \hat{\pi}_j] = i\delta_{kj}$, would be properly tractable without the need of distributions. In fact, it is in this form that we will find them later.

D.2. Problem II: The Measure

As we have mentioned several times, the main problem with an L^2 space over an infinite dimensional configuration space $Q = \mathcal{H}$ is that no obvious notion of ‘‘volume’’ exists in infinite dimensional spaces, while the Lebesgue measure gave us a unique volume notion in \mathbb{R}^n . A notion of volume is first of all well-defined in all open sets. Hence, we demand it to be a Borel measure (a measure over the Borel σ -algebra of the topological space \mathcal{H}). Now, the key of what defines ‘‘volume’’ is that it is a translation-invariant measure of sets, meaning that if we take a (measurable) set $B \subset \mathcal{H}$ (representing an ensemble of field configurations), $Volume(B) = Volume(B + \eta) \forall \eta \in \mathcal{H}$.

Proposition 2. (Eldredge, 2016) Let W be an infinite dimensional separable Banach space. There exists no translation invariant measure μ on the Borel σ -algebra of W that assigns positive and finite measure to open balls. In particular, any translation invariant Borel measure μ is either the zero measure or assigns $+\infty$ to every open set. ♦

Proof. The Key problem: By the Riesz lemma, in any ball $B(x, r)$ with $x \in W$, $r > 0$, \exists a radius $s \in (0, r)$ for which we can find countable infinity many disjoint subset balls $B(x_j, s) \subset B(x, r)$, $j \in \mathbb{N}$.

Let μ be a translation invariant Borel measure on W . By translation invariance, all $B(x_j, s)$ must have the same measure.

- If $\mu(B(x_j, s)) = \alpha > 0$, by σ -additivity

$$\mu\left(\bigsqcup_{j \in \mathbb{N}} B(x_j, s)\right) = \lim_{M \rightarrow +\infty} \sum_{j=1}^M \mu(B(x_j, s)) = \lim_{M \rightarrow +\infty} \sum_{j=1}^M \alpha = +\infty.$$

Since $\bigsqcup_{j \in \mathbb{N}} B(x_j, s) \subseteq B(x, r)$ this implies that $\mu(B(x, r)) = +\infty$. Since the ball $B(x, r)$ was arbitrary and any open set has at least a ball contained in it, μ assigns measure $+\infty$ to all open sets.

- If $\mu(B(x_j, s)) = 0$, since W is separable, it can be covered by countably many balls of radius s and therefore $\mu(W) = 0$. Hence, μ is the zero measure. *o.ε.δ.*

So, there is no volume indeed. But why did we want it to begin with? Was it really crucial?

In finite dimensional \mathbb{R}^n , a notion of translation-invariant measure is interesting because it singles out a measure among the zoo of ways to assign measures to sets: there exists only one translation-invariant Borel measure that assigns measure (volume) 1 to the unit cube $(0, 1) \times \dots \times (0, 1)$: the Lebesgue measure. However, despite this being a nice feature, it does not arrive to the heart of why we introduced measures in our discussion. In reality, we introduced measures in this context because we wanted each quantum state of Q to have an associated **probability measure**,²⁶ giving probabilities to subsets of Q . Then, the straightforward way to relate

²⁶Because experimentally we know that a quantum system is epistemologically stochastic.

probability measures to each other was by having a “background measure” $d\mu$ on Q , such that by only providing a **density** function ρ over Q , we could get very different measures via $\rho d\mu$. The $L^2(Q, \mathbb{C}, d\mu)$ space then entered the game because we needed the densities to be normalized (to yield probability measures) and because we found that $|\psi|^2$ for complex valued functions ψ over Q (obeying some particular PDE) gave the right measure of probability. Now, we chose $d\mu$ to be the Lebesgue measure, but there were more choices of background measure that would yield the same predictions (and do have an infinite dimensional version). Let us explain.

Proposition 3. For any measure space $(Q, \Sigma, d\mu)$, $L^2(Q, \mathbb{C}, d\mu)$ is a Hilbert space (Theorem 6.6 in [Folland \(1999\)](#)). Moreover (by Theorem 4.13 in [Brezis \(2011\)](#)), $L^2(Q, \mathbb{C}, d\mu)$ is separable if Q is a second countable metrizable topological space (equivalently, a separable metric space —by Theorem 16.11 in [Willard \(2012\)](#)—), with Σ being its Borel σ -algebra and if $d\mu$ a σ -finite²⁷ Borel measure. \blacklozenge

So, given two σ -finite Borel measures $d\mu_1, d\mu_2$ over $Q = \mathcal{H}$, both $L^2(Q, \mathbb{C}, d\mu_2)$ and $L^2(Q, \mathbb{C}, d\mu_1)$ are separable Hilbert spaces and hence they are unitarily equivalent. Now, if in addition, they are **mutually absolutely continuous** measures, meaning that they agree on who are the null-sets, then (by the Radon-Nikodym theorem) there exists a function $\rho^2 : Q \rightarrow \mathbb{R}$ (that is a.e. positive with respect to both μ_1 and μ_2) satisfying $d\mu_2 = \rho^2 d\mu_1$ and $d\mu_1 = \frac{1}{\rho^2} d\mu_2$. If so, the map $U : L^2(Q, \mathbb{C}, d\mu_2) \rightarrow L^2(Q, \mathbb{C}, d\mu_1)$ given by $U(\psi) := \rho\psi$ can be found to be a unitary map. And here is the catch: all pairs of wavefunctions related by the unitary yield the same probability measure on Q ! To see why:

$$\begin{aligned} \text{Prob}(q \in B \mid \psi \in L^2(Q, \mathbb{C}, d\mu_2)) &:= \int_B |\psi|^2(q) d\mu_2 = \int_B |\psi|^2(q) \rho^2(q) d\mu_1 = \\ &= \int_B (U\psi)(q) d\mu_1 =: \text{Prob}(q \in B \mid U\psi \in L^2(Q, \mathbb{C}, d\mu_1)), \end{aligned} \quad (64)$$

i.e., as measures, $|U\psi|^2 d\mu_1 = |\psi|^2 d\mu_2$.

This means that in \mathbb{R}^n if we had chosen as background measure instead of the Lebesgue measure $d^n x$, any other measure $d\mu$ mutually absolutely continuous with respect to $d^n x$, we would have obtained exactly the same quantum theory employing $L^2(\mathbb{R}^n, \mathbb{C}, d\mu)$ than the usual $L^2(\mathbb{R}^n, \mathbb{C}, d^n x)$. By “the same quantum theory” we mean that, given an arbitrary $\psi \in L^2(\mathbb{R}^n, \mathbb{C}, d\mu)$ and its identified vector $U\psi = \rho\psi \in L^2(\mathbb{R}^n, \mathbb{C}, d^n x)$:

1. The same Born rule would be given by both theories: (64) implies that ψ and $U\psi$ would attribute the same probabilities to subsets of \mathbb{R}^n (i.e., although expressed in terms of different background measures, they give rise to the same measure, $|\psi|^2 d\mu = |U\psi|^2 d^n x$)
2. The time evolution operator $\tilde{U}^t = e^{-itH}$ in $L^2(\mathbb{R}^n, \mathbb{C}, d^n x)$ would correspond to the map $U^t := U^{-1}\tilde{U}^t U$ in $L^2(\mathbb{R}^n, \mathbb{C}, d\mu)$, which is still unitary, implying (by Stone’s theorem) there exists a unique self-adjoint generator K such that $U^t = e^{-itK}$. Therefore, despite the Hamiltonians H, K would look different, the dynamics in both theories would be given by a Schrödinger equation, and U would identify solutions of one SE with solutions of the other SE.
3. If a velocity field has equivariant flow-lines for the time dependent measure generated by $t \mapsto U^t\psi$, the **same** velocity field will have equivariant flow-lines for the respective solution in the other theory $t \mapsto \tilde{U}^t(U\psi)$. this is because identified wavefunctions yield the same measure at all times:

$$|U_t\psi|^2 d\mu = |\psi_t|^2 \rho^2 d^n x = |\rho\psi_t|^2 d^n x = |U\psi_t|^2 d^n x = |UU^t U^{-1}U\psi|^2 d^n x = |\tilde{U}^t(U\psi)|^2 d^n x. \quad (65)$$

So in fact, they yield the same pilot wave theory! Therefore, both the experimental outcomes and the ontology in one and the other theory coincide and can genuinely be seen as a change of representation for the wavefunction.

This then proves that, other than due to its unique characterization, there is no compelling reason to use the Lebesgue measure as background measure for the canonical quantization. Instead, any other measure in the equivalence relation defined by mutual absolute continuity with the Lebesgue measure would have worked equivalently fine.

For later reference, let us check that in a sense, the unitary U identifying the theories also identifies the “algebras”, by checking that it is a change of representation for the canonical commutation relation:

²⁷ Q has a countable cover by finite measure sets.

4. If \hat{x}_j, \hat{p}_j are the usual position and momentum operators, with domains $D(\hat{x}_j), D(\hat{p}_j)$ in $L^2(\mathbb{R}^n, \mathbb{C}, d^n x)$, then the corresponding operators in $L^2(\mathbb{R}^n, \mathbb{C}, d\mu)$ will be the operator $U^{-1}\hat{x}_jU$ and $U^{-1}\hat{p}_jU$ on $U^{-1}D(\hat{x}_j)$ and $U^{-1}D(\hat{p}_j)$, respectively. Explicitly, $\forall \psi \in U^{-1}D(\hat{x}_j)$, it turns out that $U^{-1}\hat{x}_jU\psi = \hat{x}_j\psi$, so formally the position operator is mapped to “itself”. On the other hand, $\forall \psi \in U^{-1}D(\hat{p}_j)$, it turns out that $U^{-1}\hat{p}_jU\psi = (-i\hbar \frac{\partial \log(\rho)}{\partial x_j} + \hat{p}_j)\psi$, so the momentum operator gets an additional multiplication operator when using another background measure. Indeed, they still satisfy the commutation relations.

D.2.1 The Key Strategy

Take the Lebesgue measure dx_j on \mathbb{R} , so that $\prod_{j=1}^N dx_j = d^n x$ is the n dimensional Lebesgue measure. The limit $\lim_{N \rightarrow +\infty} \prod_{j=1}^N dx_j$ is not a measure. But:

Proposition 4. Let μ_j with $j \in \mathbb{N}$ be (inner regular²⁸) Borel **probability** measures on some metric spaces W_j . Then, by Kolmogorov’s extension theorem (see Theorems 2.4.3 and 2.4.4 in Tao (2011)), there is a unique Borel measure μ^∞ on $\prod_{j=1}^\infty W_j$ (with the product topology) such that μ^∞ is the product measure $\mu^\infty := \prod_{j=1}^\infty \mu_j$, i.e.,

- for any given finite $I \subset \mathbb{N}$ and any Borel measurable set $B_j \subset W_j$ with $j \in I$, defining the cylinder set $B := \prod_{j \in I} B_j \times \prod_{j \notin I} W_j$, we have $\mu^\infty(B) = \prod_{j \in I} \mu_j(B_j)$. More generally, if $\mu^{(I)}$ is the finite product measure of $\{\mu_j\}_{j \in I}$, for **any** Borel set of $\prod_{j \in I} W_j$ we get $\mu^\infty(B \times \prod_{j \notin I} W_j) = \mu^{(I)}(B)$.

Moreover, μ is still a probability measure on $\prod_{j=1}^\infty W_j$. ◆

So, if we had taken any probability measure in the equivalence class of measures of the Lebesgue measure dx , say, $d\mu_j$, then $\prod_{j=1}^n d\mu_j =: d^n \mu$ would still be in the same class as the n dimensional Lebesgue measure $d^n x$ (meaning they would yield equivalent pilot wave theories), but unlike with the Lebesgue measure, the infinite product measure: $\lim_{N \rightarrow +\infty} \prod_{j=1}^N d\mu$ would now have a limit product measure on $\mathbb{R}^\infty := \mathbb{R}^{|\mathbb{N}|} = \prod_{j \in \mathbb{N}} \mathbb{R}$!

Then, given the ONB coefficient space $\ell^2(\mathbb{N})$ is a subset of \mathbb{R}^∞ , we now recover the hope to find a limit theory for the tower of cut-off QFTs of §D.1 (a theory such that every cut-off theory is indeed an approximation of it). Fix a probability measure $d\mu$ on \mathbb{R} that is equivalent to the Lebesgue measure dx . For any truncation $n \in \mathbb{N}$, an equivalent cut-off theory results if we use as “background measure” $d^n x$ or $d^n \mu$, i.e., using $L^2(\mathbb{R}^n, \mathbb{C}, d^n x)$ or $L^2(\mathbb{R}^n, \mathbb{C}, d^n \mu)$. The unitary $U\psi = \rho^2\psi$ would identify the field operators as in point 4 in the last subsection. Now, while the limit space (taking $N \rightarrow +\infty$) does not exist for dx , it does for $d\mu$. Then, we would just need to find which are the operators in $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty \mu)$ such that when restricted to the cut-off spaces, they yield the field operators of the cut-off theories. Ideally, they will still satisfy the canonical commutation relation in a dense subset. Those shall be the field operators.

There are two loose ends in this strategy:

- Among the probability measures that are equivalent to the Lebesgue measure, which one gives the correct limit space $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty \mu)$? This loose end would not emerge in the case that each choice of $d\mu$ yielded an equivalent infinite product measure and hence an equivalent wavefunctional space $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty \mu)$. This issue will be solved in future work.
- How to find more explicitly the field operators in the limit space $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty \mu)$? A possible idea would be to check whether the set formed by wavefunctionals only taking non-zero values on truncated sequences $(\alpha_1, \dots, \alpha_N, 0, 0, \dots)$ (i.e., the cut-off wavefunctional set for which we know the field operators by point 4 above) is dense in $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty \mu)$, and then define $(\hat{\alpha}_j \Psi)(\alpha) := \alpha_j \Psi(\alpha)$ in a suitable dense subset of it. On the other hand, define $(\hat{\pi}_j \psi)(\alpha) := (-i\hbar \frac{\partial}{\partial \alpha_j} \log(\rho) - i\hbar \frac{\partial}{\partial \alpha_j})\psi(\alpha)$ also in some suitable dense subset (by point 4 above). Such an idea will be used in §D.4.

A heuristic that could help in the search for (a) —which seems to be the most underdetermined loose end— is that if in any “finite dimensional approximation” of configuration-space Q , call it Q_n , we find an a.e. positive function $\rho_n^2 := \otimes_{j=1}^n \rho^2 : Q_n \rightarrow \mathbb{R}$ that appears next to $d^n x$ when taking integrals of $L^2(Q_n, \mathbb{C}, d^n x)$, then, we could put together $\rho_n^2 d^n x = \prod_{j=1}^n (\rho^2 dx)$ and defining $d\mu := \rho^2 dx$, build an equivalent pilot wave theory on $L^2(Q_n, \mathbb{C}, \prod_{j=1}^n d\mu)$, where now the limit space $L^2(Q, \mathbb{C}, \prod_{j=1}^\infty d\mu)$ exists. To clarify the use of this heuristic, you can now proceed to Appendix D where we find how it can be employed for instance to derive the Feynman-

²⁸For our cases of interest, e.g., probability measures on \mathbb{R} , it will always be so.

Kac formula. In §E, our procedure will be to find an already recognized way to make QFT rigorous (using algebraic QFT and Fock spaces) and then find a representation equivalence to some $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$, with $d\mu$ chosen following the “trick” just mentioned. But before that, we must acknowledge an important detail.

D.2.2. The Price to Be Paid

To achieve a limit measure in §D.2.1, a potentially tragic sacrifice was made. The product vector space in Proposition 4, $\prod_{j \in \mathbb{N}} W_j$, which in our case of interest is \mathbb{R}^∞ , needs to be taken in the product topology. But this topology is not normable (no norm exists such that its ball topology matches the product topology).²⁹ Hence, although there will exist a well-defined separable³⁰ wavefunctional Hilbert space $L^2(\mathbb{R}^\infty, \mathbb{C}, \prod_{j=1}^\infty d\mu_j)$, the configuration space \mathbb{R}^∞ will have no norm. And a norm was key for defining differentiation of wavefunctionals with respect to fields of the configuration-space.

Remarkably, it is not clear that we needed it either, since time derivatives for paths in $L^2(\mathbb{R}^\infty, \mathbb{C}, \prod_{j=1}^\infty d\mu_j)$ are very well-defined. As such, a SCOPUG on $L^2(\mathbb{R}^\infty, \mathbb{C}, \prod_{j=1}^\infty d\mu_j)$, say $U_t = e^{itH}$, still makes sense and satisfies a SE $i\hbar \frac{d}{dt} \psi_t = H\psi_t$ for $\psi_t = U\psi_0$ and $\psi_0 \in D(H)$. Likewise, a continuity equation can still be extracted from such a SE. Also, since $L^2(\mathbb{R}^\infty, \mathbb{C}, \prod_{j=1}^\infty d\mu_j)$ itself is a Hilbert space, then we can still do calculus for operators and more general functions over it.

Moreover, we really only used to take derivatives of wavefunctions with respect to configuration-space in dense subspaces (like some $H^1(\mathbb{R}^{3N})$ in the case of $L^2(\mathbb{R}^{3N})$). Hence, we would have enough with dense subspaces that do have a norm, where taking functional derivatives does make sense. Namely, we can likely find an infinite dimensional subspace of \mathbb{R}^∞ , ideally $\ell^2(\mathbb{N})$, that does have a norm and the norm topology coincides with the subspace topology induced by \mathbb{R}^∞ —or at least that the Borel σ -algebras coincide. If so, the Borel measure on \mathbb{R}^∞ would restrict to a Borel measure for the NVS $\ell^2(\mathbb{N})$ and then, if $\mu(\ell^2(\mathbb{N})) = 1$, we could take $L^2(\ell^2(\mathbb{N}), \mathbb{C}, \mu|_{\ell^2})$ altogether as our wavefunctional space.

Other alternative tricks exist too. For example, although the product topology is not normable, there exists a distance function, namely,

$$d((\alpha_j), (\beta_j)) := \sum_{k=0}^{\infty} \frac{|\alpha_k - \beta_k|}{2^k(2 + |\alpha_k - \beta_k|)} \quad (\alpha_j)_j, (\beta_j)_j \in \prod_{j \in \mathbb{N}} W_j, \quad (66)$$

so at least it is a metric space. Then, generalized notions to those desired from calculus might be found.

Alternatively, we could accept a non-Borel measure for the probabilistic interpretation (which in the finite dimensional restrictions is nevertheless Borel). By this we mean that one could norm \mathbb{R}^∞ , allowing calculus over the norm topology and only use the measure and the product topology’s σ -algebra for the probabilistic interpretation.

For the present work, we will not need to care about this detail any-more. Just note that we will work on \mathbb{R}^∞ as a legitimate field configuration-space, assuming it contains $\ell^2(\mathbb{N})$ (the ONB coefficient space) in a suitable way.

²⁹*Proof.* Assume such a norm $\|\cdot\| : W_j \rightarrow [0, +\infty)$ existed. Then, because the unit ball B_1 around the zero vector $(0, 0, 0, \dots) \in \prod_{j \in \mathbb{N}} W_j$ would be open and the product topology is generated by cylinder sets, there would exist a cylinder subset of the unit ball, say $O \subseteq B_1$ such that $O = \prod_{j \notin I} W_j \times \prod_{j \in I} B_j$ with $I \subseteq \mathbb{N}$ finite and $B_j \subseteq W_j$ open. But if so, for any $k \in \mathbb{N} \setminus I$ and any $x_k \in W_k \setminus \{0\}$, the point $x := (0, \dots, 0, x_k, 0, \dots) \in O$. Since W_k is a vector space, $tx \in O \subseteq B_1$ for all $t > 0$, but using the homogeneity of the norm, $\|tx\| = t\|x\|$. Hence, the point $y := 2x/\|x\|$ is in the unit ball and simultaneously $\|y\| = 2$. Absurd! o.ε.δ.

³⁰By Theorem 16.4.c of Willard (2012), the product topology of separable topological spaces is separable if and only if there are at most $|\mathbb{R}|$ factors. In particular, \mathbb{R}^∞ is separable. By the distance function (66), \mathbb{R}^∞ is metrizable. Then, by Theorem 16.11 in Willard (2012), any separable metric space is second countable, thus \mathbb{R}^∞ is second countable. Finally, applying Proposition 3, we get that $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$ is separable.

D.3. Algebraic QFT and How to Use it to Build a Distinguished QFT over Field Configurations

Consider the well-known QM over finite dimensional configuration spaces. Following the narrative of Wallace (2006), one way to specify the theory would be to give a Hilbert space $(\mathcal{H}, \langle \cdot, \cdot \rangle)$ and a set of self-adjoint operators on it $\{\hat{O}_j\}_j$ for which we give names manifesting their “experimental” correlatives: e.g., “position of j -th degree of freedom” or “canonical momentum of k -th degree of freedom”. We declare that points $\psi \in \mathcal{H}$ describe laboratory preparations such that $\langle \psi, \hat{O}_j \psi \rangle$ yields the expected value for the respective “experimental correlative”. Then, we would specify a function of those operators \hat{H} , call it “Hamiltonian”, say it generates the time evolution of points ψ in the Hilbert-space as $U_t = e^{-it\hat{H}}$ and make each operator convey the effect of time evolution: $\hat{O}_j(t) := U_t \hat{O}_j U_t^{-1}$. According to the orthodox view, the set of time evolved operators contains all the predictions of the theory, so we could now forget about the U_t . Providing only the action of the operators in time, $\hat{O}_j(t)$, on \mathcal{H} together with their experimental correlatives and their algebraic relations, we would have all orthodox QM needs. Now, what if we also forgot about their action on \mathcal{H} and the fact they are operators at all, merely keeping the abstract algebraic relations between the operators (or more particularly, between the bounded functions built from those position and momentum operators), keeping as labels their experimental correlatives? Would we be able to recover the theory back? Remarkably, the answer is that for finitely many position-momentum pairs yes! For a follower of Heisenberg’s theory, this is not merely a causality, but it is the true way to approach QM. Let us introduce their approach, so that we can see how they already found a way to rigorify our driving example QFT in a rather different space than the wavefunctionals leading to field ontology. The idea is to use it as a solid scaffold or mirror, relative to which we can build our wavefunctional rigorization.

Definition 2. We say that a complex Banach space $(A, \|\cdot\|)$ (complete NVS over \mathbb{C}) equipped with an associative bilinear product \cdot such that $\|x \cdot y\| \leq \|x\| \cdot \|y\| \forall x, y \in A$ is a **Banach algebra**. \blacklozenge

Definition 3. A C^* -**algebra** is a Banach algebra $(A, \|\cdot\|, \cdot)$ equipped with an involution (a bijection that is its own inverse), $*$: $A \rightarrow A$, that has the properties of the adjoint (i-iii) and is compatible with the norm (iv):

$$(i) (x + y)^* = x^* + y^* \quad (ii) (xy)^* = y^* x^* \quad (iii) (\lambda x)^* = \bar{\lambda} x^* \quad (iv) \|xx^*\| = \|x\| \cdot \|x\| = \|x\|^2$$

for all $x, y \in A, \lambda \in \mathbb{C}$. \blacklozenge

Definition 4. We define a **representation** of the C^* -algebra $(A, \|\cdot\|, \cdot, *, \cdot)$ as a linear map $\mathcal{r} : A \rightarrow \mathcal{L}(\mathcal{H})$ (only bounded operators!) with \mathcal{H} a Hilbert space, such that

$$\mathcal{r}(a_1 \cdot a_2) = \mathcal{r}(a_1) \cdot \mathcal{r}(a_2) \quad \text{and} \quad \mathcal{r}(a^*) = \mathcal{r}(a)^* \quad \forall a_1, a_2, a \in A.$$

We say it is an **irreducible representation** when \mathcal{H} does not split in orthogonal subspaces preserved by the action of all operators in $\mathcal{r}(A)$. (It is irreducible if and only if there exists a cyclic vector for the represented operators, i.e., a vector $\psi \in \mathcal{H}$ whose image under the operators $\{\mathcal{r}(a)\psi \mid a \in A\}$ is dense). \blacklozenge

Definition 5. We say that two representations $(\mathcal{r}, \mathcal{H}), (\mathcal{r}', \mathcal{H}')$ of the same C^* -algebra A are **equivalent** if there exists a unitary isomorphism $U : \mathcal{H} \rightarrow \mathcal{H}'$ such that $\mathcal{r}(a) = U^{-1} \mathcal{r}'(a) U$ for all $a \in A$. \blacklozenge

To put it into context, assume we take a quantum theory and do the procedure we mentioned before of forgetting everything except the bounded operators and their relations. We build an abstract C^* -algebra out of them, so that the quantum theory is an irreducible representation of the abstracted C^* -algebra. If we are able to prove that there exists a unique representation of the C^* -algebra up to unitary equivalence, then, whichever representation we find of it, we will have essentially recovered the original quantum theory (as long as we remember the experimental correlatives of the operators!). In such a case, the algebra would characterize fully the quantum theory.

For finitely many quantized degrees of freedom (“finitely many position and momentum operators”) this is precisely the case by the Stone-von Neumann theorem (Reed and Simon, 1981). So, the algebraic and “Schrödinger way” of specifying the theory match in finite dimensions. However, for infinitely many quantized degrees of freedom, and hence in particular for QFT (the field operators being the “infinitely many position-momentum operators”), the theorem fails: for a C^* -algebra abstracted in such a context, there even exist uncountably many inequivalent representations. The algebraic structure alone is not enough to specify a QFT.

The reason why an algebraist still sees this abstraction technique as the foundational truth is that for them, the core of the canonical quantization process is the declaration of some “position” and “momentum”-like operators (say, the field operators) that satisfy a fixed commutation relation: hence, they claim that the key of canonical quantization is that it imposes algebraic constraints on the generated theory’s operators, independently of the Hilbert space in which we implement them. Believing this to its last consequences, they have built a very strong programme with the goal to rigorify QFT known as Algebraic QFT (AQFT) —see (Baez et al., 2014) for a (very) technical introduction.

Following their *modus operandi*, our strategy will be to abstract the algebraic relations between the heuristic field operators we developed in §B, find their corresponding C^* -algebra; there we will find a particular representation that AQFT has used to rigorify our KG QFT, even if it happens very far away from wavefunctionals: it uses the so-called, Fock spaces. This representation is legitimized because it is actually the closest one to what experimentalists use to check their cross-section experiments (they use heuristic Fock spaces to describe the asymptotically free states that are detected in the walls of particle colliders). It is completely alien to the wavefunctional approach we are suggesting, but we will find that each discretization of this Fock space QFT has an obvious equivalent wavefunctional cut-off representation. Then, using our trick in §D.2.1 to take the “no cut-off” limit, we will find an equivalent representation of the infinite dimensional C^* -algebra, but now in wavefunctional space. That is how we are going to tie the loose ends (a) and (b) to an already recognized rigorization!

D.3.1. The C^* -algebras that Abstract the Commutation Relations

This section follows from Bachmann (2017). If we take a manual on C^* -algebras, the first thing we will find in our interest’s line is the Canonical **Anti**-commutation Relation (CAR), which is

Definition 6. Given a Hilbert space \mathcal{H} , we define the CAR algebra $\mathcal{A}_-(\mathcal{H})$ as the C^* -algebra generated by the identity Id and the elements $a(f)$ with $f \in \mathcal{H}$ such that for all $f, g \in \mathcal{H}$

- $f \mapsto a(f)$ is anti-linear.
- $\{a(f), a(g)\} = 0 = \{a(f)^*, a(g)^*\}$.
- $\{a(f), a(f)^*\} = \langle f, g \rangle Id$.

This is the algebra of creation and annihilation operators abstracted from a fermionic theory. It is particularly nice because it can be represented by bounded self-adjoint operators. Yet, it is not the commutation relation we are looking for.

We would then leaf through the manual, looking for a similar thing but using commutators, something like the C^* -algebra for bosonic creation and annihilation operators, perhaps called Canonical Commutation Relations (CCR). Instead of a C^* -algebra we would find something in the line of

Definition 7. Given a Hilbert space \mathcal{H} and a set of operators $a(f)$ with $f \in \mathcal{H}$ such that for $f, g \in \mathcal{H}$ satisfy

- $f \mapsto a(f)$ is anti-linear.
- $[a(f), a(g)] = 0 = [a(f), a(g)]$
- $[a(f), a(g)^*] = \langle f, g \rangle Id$

understood as equalities over **dense domains**, we say that they satisfy the **CCR**.

We will find that the problem is that such operators cannot form a proper C^* -algebra because they will need to be unbounded operators! Instead, the clever AQFT community has found an “almost equivalent” C^* -algebra, which are essentially the exponentials (by the functional calculus) of such creation and annihilation operators. By the spectral functional calculus, the exponentials are bounded (unitary) operators and hence can indeed form a C^* -algebra. The exponential version of the CCR is called “Weyl algebra”.

Definition 8. Given a Hilbert space \mathcal{H} , we define the exponential or **Weyl CCR** algebra $\mathcal{A}_+(\mathcal{H})$ as the C^* -algebra generated by³¹ the elements $W(f)$ with $f \in \mathcal{H}$ such that for all $f, g \in \mathcal{H}$

- $W(-f) = W(f)^*$.
- $W(f)W(g) = \exp\left(\frac{-i}{2}Im\langle f, g \rangle\right)W(f + g)$.

³¹As a comment on well-definition, one can prove that if A_1 and A_2 satisfy $\mathcal{A}_+(\mathcal{H})$ there exists a unique $*$ -isomorphism $\gamma : A_1 \rightarrow A_2$ such that $a_2(f) = \gamma(a_1(f)) \forall f \in \mathcal{H}$. The same holds for $\mathcal{A}_-(\mathcal{H})$.

Note that this implies the commutation relation $W(f)W(g) = \exp(-i\text{Im}\langle g, f \rangle)W(g)W(f)$. \blacklozenge

Under usual conditions (the representation being regular³²) a representation of the Weyl CCR algebra $W(f)$ will have³³ densely defined self-adjoint generator $\hat{\varphi}(f) \forall f$ that will give bosonic creation and annihilation operators satisfying the CCR of Definition 7. But in addition it will also yield some “conjugate field operators” $\hat{\phi}(f), \hat{\pi}(f)$ satisfying the CCR we wanted to impose in the canonical quantization, $[\hat{\phi}(f), \hat{\pi}(g)] = i\langle f, g \rangle$. We will see more in a moment.

Note that the **number of quantized degrees of freedom** of a quantum theory representing the CAR or CCR (informally defined before as “number of position and momentum operators”) is the dimension of the Hilbert space \mathcal{H} that represents the C^* -algebra.

After this, the manual of C^* -algebras will show us how to build the so-called **Fock spaces** that satisfy the CAR and Weyl CCR. This then proves that at least one representation of each, the CCR and CAR, exists for any Hilbert space \mathcal{H} that we want to quantize (even for infinitely many degrees of freedom to be quantized).

Definition 9. Let there be an arbitrary complex Hilbert space \mathcal{H} .

- Define for a permutation $\sigma \in S_n$ the vector permutation map Π_σ on $\mathcal{H}^{\otimes n}$ given as the extension by linearity of $\Pi_\sigma : \psi_1 \otimes \cdots \otimes \psi_n \rightarrow \psi_{\sigma^{-1}(1)} \otimes \cdots \otimes \psi_{\sigma^{-1}(n)}$.
- For $n \in \mathbb{N}$, define $\mathcal{H}_\pm^{(n)} := \{\psi^{(n)} \in \mathcal{H}^{\otimes n} \mid \Pi_\sigma \psi^{(n)} = (\pm 1)^{\text{sign}(\sigma)} \psi^{(n)} \forall \sigma \in S_n\}$. That is, the symmetric (+) and antisymmetric (-) subspace of $\mathcal{H}^{\otimes n}$. Define also $\mathcal{H}_\pm^{(0)} := \mathbb{C}$.
- We define the **bosonic (+) and fermionic (-) Fock space** over \mathcal{H} by $\mathcal{F}_\pm(\mathcal{H}) := \bigoplus_{n=0}^{\infty} \mathcal{H}_\pm^{(n)}$. An element $\psi \in \mathcal{F}_\pm(\mathcal{H})$ is a tuple $\psi = (\psi^{(0)}, \psi^{(1)}, \dots)$ with $\psi^{(j)} \in \mathcal{H}_\pm^{(j)}$ and $\sum_{n \in \mathbb{N}} \|\psi^{(n)}\|_{\mathcal{H}}^2 < +\infty$.
- If we equip $\mathcal{F}_\pm(\mathcal{H})$ with the inner product $\langle \psi, \phi \rangle := \sum_{j=0}^{\infty} \langle \psi^{(j)}, \phi^{(j)} \rangle$, it is a Hilbert space.
- If \mathcal{H} represents the Hilbert space of some quantum theory for a single “particle”, $\mathcal{H}_\pm^{(n)}$ is the quantum theory of n identical fermionic or bosonic “particles”. In such a case, if $\|\psi\| = 1$, then $\|\psi^{(n)}\|^2$ can be interpreted as the probability that there are exactly n “particles” (for us they will be virtual particles). Hence, for such \mathcal{H} , Fock space describes the Hilbert space of a variable but finite number of particles:

$$\text{Prob}\left(\#particles \geq N \mid \psi\right) = \sum_{n \geq N} \|\psi^{(n)}\|^2 \underbrace{\rightarrow}_{N \rightarrow +\infty} 0.$$

- In analogy, whatever \mathcal{H} is, we refer to the n -th factor of the direct sum of $\mathcal{F}_\pm(\mathcal{H})$, namely $\mathcal{H}_\pm^{(n)}$, as the **“particle” sector**.
- We call $\Omega := (1, 0, 0, \dots)$ the **vacuum state** (it would represent a state with no “particle”).
- Define the set of states that are non-zero only in a finite number of sectors (describing finitely many “particles”) by $\mathcal{F}_\pm^{\text{fin}}(\mathcal{H}) := \{\psi \in \mathcal{F}_\pm(\mathcal{H}) \mid \exists N \in \mathbb{N} \text{ such that } \psi^{(n)} = 0 \forall n \geq N\}$. (It is dense in $\mathcal{F}_\pm(\mathcal{H})$!)
- We define the extension of the multiplication operator $N : \mathcal{F}_\pm^{\text{fin}}(\mathcal{H}) \rightarrow \mathcal{F}_\pm^{\text{fin}}(\mathcal{H})$ such that $N\psi = n\psi \forall \psi \in \mathcal{H}_\pm^{(n)}$, to be the **number operator**. The Fock space can be seen as the spectral diagonalization space of the number operator³⁴ because the projection to each sector is precisely its spectral PVM. The number operator is the lift³⁵ of the Id operator.

³²A representation $(\mathcal{H}, \mathcal{Z})$ of $\mathcal{A}_+(H)$ is said to be **regular** if $t \mapsto \mathcal{Z}(W(tf))$ is a SCOPUG for all $f \in \mathcal{H}$, and thus, is generated by a self-adjoint operator.

³³If the representation is regular, by Stone’s theorem, there exists a densely defined s.a. generator $\hat{\varphi}(f)$ such that $\mathcal{Z}(W(tf)) = \exp(it\hat{\varphi}(f))$ for all $f \in \mathcal{H}$. Moreover, $[\hat{\varphi}(f), \hat{\varphi}(g)] = i\text{Im}\langle f, g \rangle$ on $D(\hat{\varphi}(f)) \cap D(\hat{\varphi}(g))$. By defining $a^*(f) := 2^{-1/2}(\hat{\varphi}(f) - i\hat{\varphi}(if))$ and $a(f) := 2^{-1/2}(\hat{\varphi}(f) + i\hat{\varphi}(if))$ on $D(\hat{\varphi}(f)) \cap D(\hat{\varphi}(if))$, which is dense in \mathcal{H} , then $[a(f), a^*(g)] = \langle f, g \rangle$, satisfying Def. 7 (the CCR for creation and annihilation operators). Our desired CCR for canonical “conjugate” field operators can be obtained then as explained after equation (67).

³⁴The wavefunctional space on the other hand could be viewed as the space where the field operators are diagonal.

³⁵ Given a densely defined operator (A, D) on \mathcal{H} , we define its **lift** or **second quantization** to be the operator $d\Gamma(A)$ with domain $\Gamma(D) := \{\psi \in \mathcal{F}_\pm^{\text{fin}}(\mathcal{H}) \mid \psi^{(n)} \in \otimes_{k=1}^n D \forall n \in \mathbb{N}\}$ that acts as $d\Gamma(A)\psi^{(n)} = \sum_{j=1}^n Id \otimes \cdots \otimes A \otimes \cdots \otimes Id \psi^{(n)}$ on $\psi^{(n)} \in \mathcal{H}_\pm^{(n)}$. For example $N = \Gamma(Id)$.

- For $f \in \mathcal{H}$ define the operator $a_{\pm}(f) : \mathcal{H}^{\otimes n} \rightarrow \mathcal{H}^{\otimes(n-1)}$ removing a “particle” in state f as

$$a_{\pm}(f)(\psi_1 \otimes \cdots \otimes \psi_n) := \sqrt{n} \langle f, \psi_1 \rangle \psi_2 \otimes \cdots \otimes \psi_n$$

and $a_{\pm}(f)\mathcal{H}_{\pm}^{(0)} = 0$ (e.g., $a_{\pm}(f)\Omega = 0$). Its obvious extension $a(f) : \mathcal{F}_{\pm}(\mathcal{H}) \rightarrow \mathcal{F}_{\pm}(\mathcal{H})$ is called the **annihilation operator** of $f \in \mathcal{H}$. (As a map $f \mapsto a_{\pm}(f)$ it is anti-linear and $Na_{\pm}(f) = a_{\pm}(f)(N - Id)$).

- Its adjoint $a_{\pm}^*(f) := a_{\pm}(f)^* : \mathcal{H}^{(n-1)} \rightarrow \mathcal{H}^{(n)}$ given for $\sigma_k^{-1} := (k, 1, \dots, \hat{k}, \dots, n)$ as

$$(a_{\pm}^*(f)\psi^{(n-1)}) = \frac{1}{\sqrt{n}} \sum_{k=1}^n (\pm 1)^{k-1} \Pi_{\sigma_k} f \otimes \psi^{(n-1)}$$

is called the **creation operator** of $f \in \mathcal{H}$. Its action is exactly to take the $n - 1$ particle sector vector, to add a “particle” in the state f and then (anti)symmetrize the result. (Note $f \mapsto a_{\pm}^*(f)$ is linear). \blacklozenge

Proposition 5. Within the notation of the last definition,

- $a_{\pm}(f)$ and $a_{\pm}^*(f)$ satisfy the CAR and CCR respectively for $(-)$ and $(+)$, in particular, $[a_+(f), a_+^*(f)] = \langle f, g \rangle Id$ and $\{a_-(f), a_-^*(f)\} = \langle f, g \rangle Id$.
- Ω is a cyclic vector for the creation operators, in the sense that $\{a_{\pm}^*(f_1) \cdots a_{\pm}^*(f_n)\Omega \mid f_j \in \mathcal{H}, n \in \mathbb{N}\}$ is dense in the Fock space (any state of the Fock space can be written as a series of “creations of particles from the vacuum”). Hence, Fock space is an irreducible representation of the CAR and CCR!
- Moreover, given any “1-particle” ONB $B := \{f_j\}_{j \in \mathbb{N}} \subset \mathcal{H}$, the following is an ONB of $\mathcal{F}_+(\mathcal{H})$:

$$\left\{ \frac{1}{\sqrt{n_1! \cdots n_N!}} a_+^*(f_{j_1})^{n_1} \cdots a_+^*(f_{j_N})^{n_N} \Omega \mid N, n_k \in \mathbb{N} \text{ and } f_{j_k} \in B \text{ and } j_1 < \cdots < j_N \right\}$$

while the following is an ONB of the fermionic $\mathcal{F}_-(\mathcal{H})$

$$\{a_-^*(f_{j_1}) \cdots a_-^*(f_{j_N})\Omega \mid N \in \mathbb{N} \text{ and } f_{j_k} \in B \text{ and } j_1 < \cdots < j_N\}.$$

(Hence, we will be able to write any state in terms of the vacuum state, and this will give us the “ ρ^2 that always appears next the Lebesgue measure” that we mentioned in §D.2.1). \blacklozenge

D.3.2. A rigorous QFT over Fock-Space

Centring in the bosonic Fock space (following now the presentation by [Reed and Simon \(1975\)](#)), we will find that there is a rigorous way to arrive at the “canonically conjugate” field operators we are interested on, but in Fock space instead of field configuration-space. This is not exactly what we were expecting from the quantization recipe in B, but as we explained, it will be the perfect proxy to find the rigorous field-configuration space theory.

Definition 10. Taking the Fock space construction of Def. 9, we define the **Segal field operators** (that end up yielding what physicists call the *real* free bosonic field’s quantization) as $\hat{\varphi}_+(f) := \frac{1}{\sqrt{2}}(a_+(f) + a_+^*(f))$ for $f \in \mathcal{H}$. \blacklozenge

They also satisfy the CCR in $\mathcal{F}_+^{fin}(\mathcal{H})$,³⁶ but now in a shape that is closer to field operators (note the i):

$$[\hat{\varphi}_+(f), \hat{\varphi}_+(g)] = i \operatorname{Im} \langle f, g \rangle. \quad (67)$$

Now, the main difference between $\hat{\varphi}_+$ and our desired field operators $\hat{\phi}, \hat{\pi}$ is that in general, the Segal field operator takes any field in the “1-particle” Hilbert space \mathcal{H} , where \mathcal{H} is taken to be a **complex** Hilbert space in

- Given a unitary U on \mathcal{H} , we call the unitary $\Gamma(U)$ on $\mathcal{F}_{\pm}(\mathcal{H})$ such that it acts as $\otimes_{k=1}^n U$ on $\mathcal{H}_{\pm}^{(n)}$ the **lifted** or **second quantized propagator**. If e^{itA} is a SCOPUG on \mathcal{H} then $\Gamma(e^{itA}) = e^{itd\Gamma(A)}$ is a SCOPUG on $\mathcal{F}_{\pm}(\mathcal{H})$. Hence, there is a direct way to implement free particle dynamics in the Fock space.

³⁶In particular, they are symmetric on $\mathcal{F}_+^{fin}(\mathcal{H})$ and extend to self-adjoint operators. Hence, they generate unitaries $W_+(f) := \exp(i\hat{\varphi}(f))$ on the whole $\mathcal{F}_+(\mathcal{H})$ which are a representation of the Weyl algebra (the exponential CCR).

the present construction,³⁷ (e.g., $H^2(\mathbb{R}^3, \mathbb{C}, d^3x)$ for our driving example). Instead, we wanted to build a QFT over the configuration space of **real** valued fields (say, $H^2(\mathbb{R}^3, \mathbb{R}, d^3x)$). Hence, we need to “uncomplexify” \mathcal{H} , taking its underlying real Hilbert space (chosen by some conjugation, which in our case is) given by complex conjugation.

Definition 11. Let \mathcal{H} be a complex Hilbert space and C a conjugation (a conjugate-linear, isometric involution, e.g., the point-wise complex conjugation in $L^2(\mathbb{R}^n)$). Define $\mathcal{H}_R := \{f \in \mathcal{H} \mid Cf = f\}$ (which for the complex conjugation yields the subspace of real valued functions). In particular, we get that $\mathcal{H} = \mathcal{H}_R \oplus i\mathcal{H}_R$. Then, for $f \in \mathcal{H}_R$, we define $\hat{\phi}(f) := \hat{\varphi}_+(f)$ and $\hat{\pi}(f) := \hat{\varphi}_+(if)$ as the **canonical conjugate field operators**. \blacklozenge

Proposition 6. Given the notation of the last definition,

- $\forall f \in \mathcal{H}_R$ the operators $\hat{\phi}(f), \hat{\pi}(f)$ are essentially self-adjoint on \mathcal{F}_+^{fin} .
- $\{\hat{\phi}(f) \mid f \in \mathcal{H}_R\}$ is a commuting family of self-adjoint operators (likewise for $\hat{\pi}(\cdot)$).
- Ω is cyclic for $\{\hat{\phi}(f)\}$. In particular, given an ONB of \mathcal{H} such that $\mathcal{B} := \{f_j\}_{j \in \mathbb{N}} \subset \mathcal{H}_R$, the span of

$$\left\{ P\left(\hat{\phi}(f_{k_1}), \dots, \hat{\phi}(f_{k_N})\right)\Omega \mid N \in \mathbb{N}; f_k \in \mathcal{B}; P \text{ a complex valued polynomial in } N \text{ real variables} \right\}$$

is dense in \mathcal{H} .

- For all $f, g \in \mathcal{H}_R$, the CCR we were looking for, $[\hat{\phi}(f), \hat{\pi}(g)] = i \langle f, g \rangle$, holds on the dense set \mathcal{F}^{fin} .
- For all $f, g \in \mathcal{H}_R$, also the Weyl CCR holds, $e^{i\hat{\phi}(f)}e^{i\hat{\pi}(g)} = e^{-i\langle f, g \rangle}e^{i\hat{\pi}(g)}e^{i\hat{\phi}(f)}$. \blacklozenge

If we did not care about field configuration-space wavefunctionals, then, for $\mathcal{H} = H^2(\mathbb{R}^3, \mathbb{C}, d^3x)$, the Fock space $\mathcal{F}_+(\mathcal{H})$ with the field operators $\hat{\phi}(f)$ and $\hat{\pi}(f)$ and $f \in H^2(\mathbb{R}^3, \mathbb{R}, d\mu)$ would already satisfy most of the other conditions of canonical quantization given heuristically in §B. In fact, as we mentioned already, it is the standard way to make it mathematically rigorous and physicists employ a heuristic version of this construction to define asymptotic states of their experiments. However, we lost two main things in skipping wavefunctional spaces: (i) how simple it was to obtain the Hamiltonian for second quantized dynamics, and (ii) the opportunity to have a pilot wave theory with field ontology behind it, which is our leitmotif in this document. Let us recover them.

D.3.3. A Baby Version of the Grand Finale

What we are going to do here for one degree of freedom (find an equivalence between the Schrödinger and Fock representations) is exactly what gives the wavefunctional picture if done for infinitely many degrees of freedom. The Schrödinger picture will be exactly the wavefunctional picture.

Quantum mechanics for a single degree of freedom/single particle is usually given in the so-called Schrödinger representation, which is a representation of the $\mathcal{A}_+(\mathbb{C})$ Weyl CCR algebra on $L^2(\mathbb{R}, \mathbb{C}, dx)$ (so, not Fock space) given, as explained by [Bachmann \(2017\)](#), via

$$z \left(W(s + it) \right) := e^{\frac{i}{2}st} U(s) V(t) \quad \text{where } \forall \psi \in L^2(\mathbb{R}) \begin{cases} (U(s)\psi)(x) = e^{is}\psi(x) & \text{(global phase)} \\ (V(t)\psi)(x) = \psi(x + t) & \text{(translation)} \end{cases}. \quad (68)$$

As SCOPUGs, by Stone’s theorem there exists a unique generator for $U(s), V(t)$, being the position operator $\hat{x} = x$ and the momentum operator $\hat{p} = -i\hbar\partial_x$, respectively. In a dense set (say Schwartz functions) they satisfy the ordinary CCRs, $[\hat{x}, \hat{p}] = i \text{Id}$ etc., as is well-known.

But now we also know that there must exist a Fock space representation of $\mathcal{A}_+(\mathbb{C})$ and it is certainly not the one we described. It is given in $\mathcal{F}_+(\mathbb{C}) := \bigoplus_{n=0}^{\infty} \mathbb{C}^{\otimes n}$. Note that $\mathbb{C}^{\otimes n}$ is 1 dimensional for all $n \in \mathbb{N}$ and has ONB $\{1 \otimes \dots \otimes 1\}$. Hence, an ONB for $\mathcal{F}_+(\mathbb{C})$ is given by $\{1, 1 \otimes 1, 1 \otimes 1 \otimes 1, \dots\}$ and a state $\psi \in \mathcal{F}_+(\mathbb{C})$

³⁷The Segal field operators on the bosonic Fock space can be used to build a rigorous quantization of the Klein-Gordon equation taking as \mathcal{H} the complex L^2 space over the hyperbola of Minkowski space representing the momentum-space mass-shell. Then, one uses field operators that take fields on Minkowski spacetime instead of a single slice as we do. But the input to the field operators can be restricted adequately so that a rigorous version of the quantization in single-time slices like the one we are aiming at, can be obtained as a corollary (see [Reed and Simon, 1975](#)).

is $\psi = (c_0, c_1 1 \otimes 1, c_2 1 \otimes 1 \otimes 1, \dots)$, so a square summable sequence of complex numbers $(c_j)_j \in \ell^2(\mathbb{N}_0, \mathbb{C})$. Almost as if they were the ONB expansion coefficients of a vector in $L^2(\mathbb{R})$. (Spoiler: they secretly are!)

So we obtained two different representations of the Weyl CCR and both can be checked to be irreducible.³⁸ Hence, because the number of quantized degrees of freedom is finite, i.e., $\dim(\mathcal{H}) = \dim(\mathbb{C}) = 1$, by the Stone-von Neumann theorem they must be equivalent representations of the CCR. We can explicitly find their unitary equivalence as follows.

For this, let us first find a ‘‘Fock space structure’’ hiding in $L^2(\mathbb{R})$. Define the operators

$$\hat{a}_S(1) := \frac{1}{\sqrt{2}}(\hat{x} + i\hat{p}) \quad \text{and} \quad \hat{a}_S^*(1) := \frac{1}{\sqrt{2}}(\hat{x} - i\hat{p}). \quad (69)$$

They also satisfy the CCR (in the same common dense domain), but now in the creation and annihilation operator version we saw in Def. 7, namely $[\hat{a}_S(1), \hat{a}_S^*(1)] = Id$ etc. Moreover, there is a vector that acts as the ‘‘vacuum’’ too: the solution to $a_S(1)\Omega_S = 0$,

$$\left(x + \frac{\partial}{\partial x}\right)\Omega_S(x) \equiv 0 \iff \Omega_S(x) \equiv \frac{1}{\pi^{1/4}}e^{-x^2/2}. \quad (70)$$

This is nothing but the ground-state of the harmonic oscillator (note there is no Hamiltonian yet!). What is more, this analogue of the vacuum is also cyclic with respect to repeated application of $a^*(1)$ (the span of $\{a^*(1)^n\Omega_S \mid n \in \mathbb{N}\}$ is dense in $L^2(\mathbb{R})$) because up to a norm factor $1/\sqrt{n!}$ they form an ONB:

$$\psi_n := a^*(1)^n\Omega_S = \frac{1}{2^{n/2}\pi^{1/4}}\left(x - \frac{\partial}{\partial x}\right)^n e^{-x^2/2} = \frac{1}{2^{n/2}\pi^{1/4}}e^{-x^2/2}H_n(x) \quad (71)$$

are the eigen-energy states of the harmonic oscillator, where $H_n(x) := e^{x^2/2}(x - \partial_x)^n e^{-x^2/2}$ is the n -th Hermite polynomial.³⁹ Because they form an ONB of $L^2(\mathbb{R})$, defining $\mathcal{H}^{(n)} := \text{span}\{\psi_n\} = \text{span}\{a_S^*(1)^n\Omega_S\}$, we get $L^2(\mathbb{R}) = \bigoplus_{n=0}^{\infty}\mathcal{H}^{(n)}$. But this is exactly a ‘‘Fock space structure’’! This suggests the unitary relating the representations:

$$U : L^2(\mathbb{R}) \rightarrow \mathcal{F}_+(\mathbb{C}) \quad \text{mapping} \quad \psi = \sum_{n=0}^{\infty} c_n \psi_n \mapsto (c_0, c_1 1 \otimes 1, c_2 1 \otimes 1 \otimes 1, \dots). \quad (72)$$

As promised, Fock space elements turned out to represent expansion coefficients! Importantly, the creation and annihilation operators (69) are mapped to those of Fock space. Interestingly, the canonical conjugate fields of Fock space are identified with the Schrödinger representation’s position and momentum operators:

$$\hat{\phi}_S(1) := \hat{\varphi}_S(1) := \frac{1}{\sqrt{2}}(\hat{a}_S(1) + \hat{a}_S^*(1)) = \hat{x} \quad \text{and} \quad \hat{\pi}_S(1) := \hat{\varphi}_S(i) := \frac{1}{\sqrt{2}}(-i\hat{a}_S(1) + i\hat{a}_S^*(1)) = \hat{p}. \quad (73)$$

Then, it should not be a surprise that for the wavefunctional picture of the grand finale, the canonically conjugate field operators of the Fock space will be exactly the ones mapped to the field operators as defined in §B.

With all, $U : L^2(\mathbb{R}) \rightarrow \mathcal{F}_+(\mathbb{C})$ maps:

$$\begin{aligned} \Omega_S = \pi^{-1/4}e^{-x^2/2} = \psi_0 &\mapsto \Omega = (1, 0, 0, 0, \dots) \\ (a_S^*(1)\Omega_S = \psi_1 &\mapsto (a_+^*(1)\Omega = (0, 1, 0, 0, \dots) \\ (a_S^*(1)^2\Omega_S = \psi_2 &\mapsto (a_+^*(1)^2\Omega = (0, 0, 1, 0, \dots) \\ (a_S^*(1)^n\Omega_S = \psi_n &\mapsto (a_+^*(1)^n\Omega = \underbrace{(0, 0, \dots, 1, 0, \dots)}_{n-1} \end{aligned} \quad (74)$$

$$\hat{N}_S := a_S^*(1)a_S(1) = \frac{1}{2}(\hat{x} - i\hat{p})(\hat{x} + i\hat{p}) \mapsto \hat{N} = a_+^*(1)a_+(1)$$

$$\begin{aligned} \hat{\phi}_S(1)\Omega_S = \pi^{-1/4} x e^{-x^2/2} &\mapsto \hat{\varphi}(1)\Omega = \frac{1}{\sqrt{2}}(a_+(1) + a_+^*(1)) = (0, 1, 0, \dots) \\ \hat{\phi}_S(1)^2\Omega_S = \pi^{-1/4} x^2 e^{-x^2/2} &\mapsto \hat{\varphi}(1)^2\Omega = (1, 0, 1, 0, \dots) \\ \hat{\pi}_S(1)\Omega_S = \pi^{-1/4} (-i\hbar\partial_x) e^{-x^2/2} &\mapsto \hat{\pi}(1)\Omega = \frac{1}{\sqrt{2}}(a_+(i) + a_+^*(i)) = (0, i, 0, \dots) \end{aligned} \quad (75)$$

³⁸Actually, looking for the obvious cyclic vectors will reveal the unitary equivalence.

³⁹Note there is no Hamiltonian in game, so we are just using the ONB with no additional meaning, but if there was, $H = -\frac{1}{2}\partial_x^2 + \frac{1}{2}x^2$ would be also equal to $H = a^*a + 1/2$.

Therefore, from the Schrödinger picture's perspective what the creation and annihilation operators “create” or “annihilate” is an excitation or relaxation on the ONB. Creating a particle in the Fock picture can be seen as taking the wavefunction of the Schrödinger picture to a higher excited eigenstate (jargon that only makes strict sense to use if we really had that the governing Hamiltonian was the harmonic oscillator's). Hence, the number of bosons accounts for the excitation number of the harmonic oscillator: having a Fock state with three bosons is nothing but having the particle of the Schrödinger picture in the third excited state. Having the vacuum state, hence no bosons, is nothing but having the Schrödinger particle in the ground state. Thus, from this perspective, the particles we were apparently counting in the Fock space are not actual particles: we call them **virtual particles**.

The same will happen in the final demonstration: there, we will find an equivalence between the Fock space of an infinite dimensional Hilbert space (instead of the 1D Hilbert space \mathbb{C}) and a Schrödinger picture (the wavefunctional picture) having infinitely many position and momentum operators. There, the Fock space particles will have the chance to occupy different states from an ONB, $\{g_k\}_k$, and each occupation of a state g_k will correspond to an excitation of a different “harmonic oscillator” (one oscillator per each of the g_k). Creating a Fock boson in the k -th mode g_k will mean an excitation of the wavefunctional to a “higher energy” along the k -th axis.

A few remarks:

- We would like to appreciate explicitly that there was no reference to any Hamiltonian in this development. The harmonic oscillator only appeared because we chose the position and momentum as the Schrödinger “Segal operators” (to be identified with those in Fock space). They happen to have as associated creation and annihilation operators, the ladder operators of the ONB of the harmonic oscillator. However, unitary maps between Hilbert spaces are not unique and neither are representation equivalences. One could take any ONB $\{f_n\}_{n \in \mathbb{N}_0} \subset L^2(\mathbb{R})$ (other than the harmonic oscillator) and define its creation and annihilation operators by: $\hat{A}f_0 := 0$ and $\hat{A}f_n := \sqrt{n}f_{n-1}$ for $n > 0$, which extend then by linearity. That is, define $\hat{A} := \sqrt{j}f_{j-1}\langle f_j, \cdot \rangle$, or in matrix form:

$$\left(\langle f_i, Af_j \rangle \right)_{ij} := \begin{pmatrix} 0 & \sqrt{1} & 0 & 0 & 0 & \dots \\ 0 & 0 & \sqrt{2} & 0 & 0 & \dots \\ 0 & 0 & 0 & \sqrt{3} & 0 & \dots \\ \vdots & \vdots & \vdots & \vdots & \vdots & \ddots \end{pmatrix}. \quad (76)$$

It is certainly an unbounded operator, but has a dense domain in the subspaces formed by the span of finitely many ONB elements (i.e., finitely many virtual particles or excitations), just like in Fock space. And again, A and A^* would satisfy the CCR in that domain, $[A, A^*] = Id$. Then, we would find that the unitary

$$U : L^2(\mathbb{R}) \rightarrow \mathcal{F}_+(\mathbb{C}) \quad \text{that maps} \quad \psi = \sum_{n=0}^{\infty} c_n f_n \mapsto (c_0, c_1 1 \otimes 1, c_2 1 \otimes 1 \otimes 1, \dots) \quad (77)$$

would allow us to use Fock space as a bookkeeper of the excitation levels and linear combinations of this other ONB. It would then be interesting to analyse what the “field operators” $\hat{\varphi}(1)$ and $\hat{\pi}(1)$ correspond to back in $L^2(\mathbb{R})$ in this correspondence, namely, what

$$\hat{\phi}'_S(1) := \hat{\varphi}'_S(1) := \frac{1}{\sqrt{2}}(A + A^*) \quad \text{and} \quad \hat{\pi}'_S(1) := \hat{\varphi}'_S(i) := \frac{1}{\sqrt{2}}(-iA + iA^*) \quad (78)$$

would correspond to in terms of \hat{x} and \hat{p} .

- To complete the picture, consider now $\mathcal{A}_+(\mathbb{C}^n)$ (the CCR for n quantized degrees of freedom). It has a Fock space representation over $\mathcal{F}_+(\mathbb{C}^n)$, which by the Stone-von Neumann theorem must be equivalent to the Schrödinger representation of $L^2(\mathbb{R}^n, \mathbb{C}, d^n x)$ describing n actual position and momentum operators. The explicit representation equivalence is essentially given by tensor products of the equivalence for $n = 1$.⁴⁰ The final trick of the manuscript is then to take the limit $n \rightarrow +\infty$ using the strategy of §D.2.1.

⁴⁰And the same thing will happen: creating/annihilating bosons in the k -th among the now n available modes, will account in the Schrödinger picture to promote or relax a tensor product ONB along the axis of the k -th mode. In our context, this Schrödinger picture will not represent particles, but instead they will be a cut-off of a mode expansion to n mode coefficients.

- Last but not least, notice that in Fock space there is a distinguished and **fixed** state called “vacuum”, while in the analogous Schrödinger picture, the “vacuum” is just a particular state like any other. What characterizes the “vacuum” in both spaces is that it coincides with the ground state of some Hamiltonian. Hence, in a setting where there are interactions, and therefore potentially time varying ground states of the Hamiltonian, perhaps working in the Schrödinger representation is more convenient, because the “vacuum” would change in time, which is conflictive if we have a fixed state called vacuum like in Fock space. This might be critical in the case of QFTs with interactions!

D.4. A rigorous QFT with configuration-space \mathbb{R}^∞ : “a Ceiling for the Infinite Tower?!”

“And the *Grand Finale*:” an attempt to construct rigorously a meaningful QFT over the configuration-space of expansion mode coefficients: $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty\mu)$ —a “ceiling” for the “infinite tower” of cut-off theories?

On the one hand, by choosing equivalent measures $d\mu$ to the Lebesgue measure dx , that unlike the Lebesgue, allow an infinite product measure $d^\infty\mu$, we found a way to build a quantum theory on \mathbb{R}^∞ , some $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty\mu)$. Moreover, such infinite product measures allowed us to take the resulting theory as the limit of discrete/finite dimensional cut-off approximations, via partial product measures $d^N\mu$. Still, we found a big number of possible such measures $d\mu$ to build our “arrow towards the ceiling”.

To remedy this, by using Fock space $\mathcal{F}_+(\mathcal{H})$, we have found a rigorous “second quantization” of the complexification of the field configuration-space \mathcal{H} . This gathered many of the features we looked for in the heuristic prescription of §B, but lacked the leitmotif: it is not built on a field configuration space like $L^2(\mathcal{H}, \mathbb{C}, d\mu)$, neither in its analogous (extended⁴¹) “expansion coefficient” space $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty\mu)$. But, if we find a unitary giving a representation equivalence between the field operators in Fock space and some field operators that act as multiplication by functions in $L^2(\mathbb{R}^\infty, \mathbb{C}, d^\infty\mu)$ for some $d\mu$, the objective of this document will be achieved.

Proposition 7. Given:

- a complex separable Hilbert space \mathcal{H} ,
- a conjugation C on \mathcal{H} with “real subspace” $\mathcal{H}_R := \{\psi \in \mathcal{H} \mid C\psi = \psi\}$ (which is the configuration-space we want to “second quantize”),
- the bosonic Fock space $\mathcal{F}_+(\mathcal{H})$ equipped with canonically conjugate field operators $\{\hat{\phi}(f)\}_{f \in \mathcal{H}_R}$ and $\{\hat{\pi}(f)\}_{f \in \mathcal{H}_R}$,
- any fixed ONB $\{f_j\}_{j \in \mathbb{N}_0}$ of \mathcal{H} that is made of elements of \mathcal{H}_R ,

there exist:

- a Borel (probability) measure $d\mu$ on $\mathbb{R}^\infty := \mathbb{R}^{|\mathbb{N}|} := \prod_{j \in \mathbb{N}} \mathbb{R}$ with product topology, giving the separable Hilbert space $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$, (roughly, the wavefunctionals over expansion coefficients),
- a unitary map $S : \mathcal{F}_+(\mathcal{H}) \rightarrow L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$,

such that,

- $\forall f \in \mathcal{H}_R$ the field operators $\hat{\phi}(f)$ are identified with multiplication operators. For the ONB $\{f_j\}_j$ in particular, $S\hat{\phi}(f_j)S^{-1} = \alpha_j$ (which are the “position operators” along the corresponding “axis” of $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$),
- the vacuum of Fock space Ω is identified with the function $\Psi_0(\alpha) \equiv 1$ in $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$, i.e., $S\Omega = \Psi_0$. ♦

Proof. Courtesy of **Reed and Simon (1975)**. We first construct the unitary for finite dimensional subspaces:

- **Step 1:** Fix $N \in \mathbb{N}$ and let $\{g_k\}_{k=1}^N$ be a finite subset of the ONB $\{f_j\}_{j \in \mathbb{N}}$. Define

$$\mathcal{F}_N := \text{closure} \left\{ P \left(\hat{\phi}(g_1), \dots, \hat{\phi}(g_N) \right) \Omega \mid P \text{ a complex polynomial of } N \text{ arguments} \right\} \subset \mathcal{F}_+(\mathcal{H}). \quad (79)$$

(The span of repeated application on the vacuum by the field operators of the first N modes of the ONB.)

⁴¹Technically, the expansion coefficient space should be $\ell^2(\mathbb{N})$, which is a subspace of the vector space \mathbb{R}^∞ (precisely, the subspace for which the ℓ^2 -norm of a sequence is finite).

- **Step 2:** By Proposition 6, $\hat{\phi}(g_k), \hat{\pi}(g_l)$ are essentially self-adjoint on $\mathcal{F}_+^{fin}(\mathcal{H}) \cap \mathcal{F}_N$ and they satisfy the Weyl relation on \mathcal{F}_N :

$$e^{i\hat{\phi}(g_k)} e^{i\hat{\pi}(g_l)} = e^{-ist\delta_{kl}} e^{is\hat{\pi}(g_l)} e^{it\hat{\phi}(g_k)}, \quad (80)$$

so they are a representation of the Weyl CCR with finitely many (N) ‘‘position and momentum operator pairs’’ $\hat{\phi}(g_k), \hat{\pi}(g_l)$. Moreover, Ω is cyclic by definition of \mathcal{F}_N , so it is an irreducible representation of the Weyl CCR. By the Stone-Von Neumann theorem there exists a unitary equivalence with the N ‘‘particle’’ Hilbert space $L^2(\mathbb{R}^N, \mathbb{C}, d^N x)$ that maps the field and momentum operators to the position and momentum operators there, i.e., a unitary isomorphism $\tilde{S}^{(N)} : \mathcal{F}_N \rightarrow L^2(\mathbb{R}^N, \mathbb{C}, d^N x)$:

$$\begin{cases} \tilde{S}^{(N)} \hat{\phi}(g_k) (\tilde{S}^{(N)})^{-1} = \alpha_k =: \hat{\alpha}_k \\ \tilde{S}^{(N)} \hat{\pi}(g_k) (\tilde{S}^{(N)})^{-1} = -i\hbar \frac{\partial}{\partial \alpha_k} =: \hat{p}_k \end{cases} \quad \text{for all } k \in \{1, \dots, N\}. \quad (81)$$

Note that we will denote the points of \mathbb{R}^∞ (and similarly of \mathbb{R}^N) by $\alpha := (\alpha_1, \alpha_2, \alpha_3, \dots)$ with $\alpha_j \in \mathbb{R}$, to make manifest our intention to use this as the space of coefficients of some expansion modes.

- **Step 3:** Since by definition 11, $\hat{\phi}(g_k) = (a_+(g_k) + a_+^*(g_k))/\sqrt{2}$ and $\hat{\pi}(g_k) = i(-a_+(g_k) + a_+^*(g_k))/\sqrt{2}$, using that $a_+(g_k)\Omega = 0$, we can verify the identity

$$\left(\hat{\phi}(g_k)^2 + \hat{\pi}(g_k)^2 - 1 \right) \Omega = 0. \quad (82)$$

Now, define Ψ_0 to be the vector identified with the vacuum of Fock space Ω , i.e., $\Psi_0 := \tilde{S}^{(N)}\Omega$. Then, plugging $\hat{\alpha}_k = (\tilde{S}^{(N)})^{-1} \hat{\phi}(g_k) \tilde{S}^{(N)}$ and $\hat{p}_k = (\tilde{S}^{(N)})^{-1} \hat{\pi}(g_k) \tilde{S}^{(N)}$ in equation (82), we find it to imply

$$\left(\frac{1}{2} \hat{p}_k^2 + \frac{1}{2} \hat{\alpha}_k^2 \right) \Psi_0 = \frac{1}{2} \Psi_0. \quad (83)$$

Adding together the equations for $k \in \{1, \dots, N\}$ we get the eigenfunction equation

$$\sum_{k=1}^N \left(\frac{1}{2} \hat{p}_k^2 + \frac{1}{2} \hat{\alpha}_k^2 \right) \Psi_0 = \frac{N}{2} \Psi_0, \quad (84)$$

which is the equation defining the ground state of N independent harmonic oscillators. It has as solution the Gaussian $\Psi_0 = \pi^{-N/4} \exp(-\sum_{k=1}^N \frac{\alpha_k^2}{2})$. Hence:

$$\tilde{S}^{(N)}\Omega = \frac{1}{\pi^{N/4}} e^{-\sum_{k=1}^N \frac{1}{2} \alpha_k^2}. \quad (85)$$

(The Fock vacuum is identified with a tensor product of N Gaussians in the cut-off wavefunctional space.)

- **Step 4:** By definition, any state in \mathcal{F}_N is a limit of polynomials of the field operators acting on Ω . By the equivalence \tilde{S} , any state in $L^2(\mathbb{R}^N, \mathbb{C}, d^N x)$ is spanned by acting with polynomial multiplication operators on the Gaussian Ψ_0 . Recalling the idea of §D.2.1 about finding a $\prod_{j=1}^N \rho^2$ that would ‘‘always’’ be present in the states besides the Lebesgue measure, we found it: ρ^2 is the Gaussian $\pi^{-1/2} \exp(-\alpha_k^2)$!

We define the measure $d\mu_k := \pi^{-1/2} e^{-x_k^2} dx_k$. It is mutually absolutely continuous with respect to the 1D Lebesgue measure dx_k and it is a probability measure, so-called Gaussian measure. Define the unitary map

$$\begin{aligned} U : L^2(\mathbb{R}^N, \mathbb{C}, d^N x) &\longrightarrow L^2\left(\mathbb{R}^N, \mathbb{C}, \prod_{k=1}^N d\mu_k\right) = L^2\left(\mathbb{R}^N, \mathbb{C}, \pi^{-N/2} e^{-\sum_{k=1}^N x_k^2} d^N x\right) \\ \eta(\alpha) &\longmapsto \pi^{N/4} e^{+\sum_{k=1}^N \frac{1}{2} \alpha_k^2} \eta(\alpha) \end{aligned} \quad (86)$$

giving an equivalent ‘‘cut-off’’ quantum theory (even an equivalent underlying pilot-wave theory), but now with a nicer background measure to take the limit. As we explained in §D.2, $L^2(\mathbb{R}^N, \mathbb{C}, d^N x)$ does not have any limit as $N \rightarrow \infty$, while $L^2(\mathbb{R}^N, \mathbb{C}, \prod_{k=1}^N d\mu_k)$ does have. Hence, the unitary we are really interested on to go from Fock space to a limit wavefunctional space is $S^{(N)} := U\tilde{S}^{(N)}$. According to this unitary,

$$\begin{aligned} S^{(N)} : \mathcal{F}_N &\longrightarrow L^2(\mathbb{R}^N, \mathbb{C}, \prod_{k=1}^N d\mu_k) \\ \eta &\longmapsto U\tilde{S}^{(N)}\eta \end{aligned} \quad (87)$$

$$\begin{cases} (S^{(N)}\Omega)(\alpha) & \equiv 1 \\ S^{(N)}\hat{\phi}(g_k)(S^{(N)})^{-1} & = \alpha_k =: \hat{\alpha}_k \\ S^{(N)}\hat{\pi}(g_k)(S^{(N)})^{-1} & = i\hbar\alpha_k - i\hbar\frac{\partial}{\partial\alpha_k} =: \tilde{p}_k \end{cases} \quad (88)$$

(The price we pay is the extra term in the momentum operator.)

- **Step 5:** We compute an identity. For any set of polynomials in one variable P_1, \dots, P_N :

$$\begin{aligned} \left\langle \Omega, P_1(\hat{\phi}(g_1)) \cdots P_N(\hat{\phi}(g_N))\Omega \right\rangle_{\mathcal{F}} & \stackrel{\substack{\equiv \\ S^{(N)} \text{ is unitary}}}{=} \left\langle 1, P_1(\alpha_1) \cdots P_N(\alpha_N)1 \right\rangle_{L^2(\mathbb{R}^N, \prod_{j=1}^N d\mu_j)} = \\ & = \int_{\mathbb{R}^N} P_1(\alpha_1) \cdots P_N(\alpha_N) \prod_{k=1}^N d\mu_k = \star^{42} \end{aligned} \quad (89)$$

- **Step 6:** Now consider $\mathbb{R}^\infty = \prod_{k \in \mathbb{N}} \mathbb{R}$ with the σ -algebra generated by countable products of Borel measurable sets of \mathbb{R} , call it \mathcal{A} . It is the same σ -algebra as the one generated by the product topology, so by Kolmogorov's extension theorem the infinite product measure $d\mu := \prod_{k=1}^\infty d\mu_k$ exists uniquely and is a probability measure on $(\mathbb{R}^\infty, \mathcal{A})$.

- **Step 7:**

- (a) The functionals making use of only N coordinates of its input points $\alpha := (\alpha_1, \alpha_2, \dots) \in \mathbb{R}^\infty$, defined by $G^{P;k_1, \dots, k_N} : \mathbb{R}^\infty \rightarrow \mathbb{C}$ with $G^{P;k_1, \dots, k_N}(\alpha) := P(\alpha_{k_1}, \dots, \alpha_{k_N})$ for fixed polynomials P over N variables, form a vector space of functionals

$$B_{poly} := \left\{ G^{P;k_1, \dots, k_N} \mid N \in \mathbb{N}; k_1, \dots, k_N \in \mathbb{N} \text{ and } P \text{ a complex polynomial of } N \text{ arguments} \right\} \quad (90)$$

that is dense in $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$.

- (b) Similarly, the space of elements $\eta \in \mathcal{F}_+(\mathcal{H})$ that can be written as $P(\hat{\phi}(f_{k_1}), \dots, \hat{\phi}(f_{k_N}))\Omega$ with $N \in \mathbb{N}$ for some complex polynomial P of N variables and some ONB elements f_{k_j} , i.e.,

$$\tilde{B}_{poly} := \left\{ P(\hat{\phi}(f_{k_1}), \dots, \hat{\phi}(f_{k_N}))\Omega \mid N \in \mathbb{N}; k_1, \dots, k_N \in \mathbb{N} \text{ and } P \text{ a complex poly. of } N \text{ arguments} \right\} \quad (91)$$

is dense in $\mathcal{F}_+(\mathcal{H})$.

- (c) Define the following linear isomorphism between both dense subsets:

$$\begin{aligned} S : \quad \tilde{B}_{poly} \subset \mathcal{F}_+(\mathcal{H}) & \longrightarrow B_{poly} \subset L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu) \\ P(\hat{\phi}(f_{k_1}), \dots, \hat{\phi}(f_{k_N}))\Omega & \longmapsto G^{P;k_1, \dots, k_N}. \end{aligned} \quad (92)$$

- **Step 8:** Given an arbitrary polynomial of N variables and order M , there are coefficients $c_{l_1, \dots, l_N} \in \mathbb{C}$ such that, $P(x_{k_1}, \dots, x_{k_N}) := \sum_{l_1, \dots, l_N=1}^M c_{l_1, \dots, l_N} x_{k_1}^{l_1} \cdots x_{k_N}^{l_N}$. Then,

$$\left\| P(\hat{\phi}(f_{k_1}), \dots, \hat{\phi}(f_{k_N}))\Omega \right\|_{\mathcal{F}_+(\mathcal{H})}^2 = \sum_{l_1, \dots, l_N=1}^M \sum_{m_1, \dots, m_N=1}^M c_{l_1, \dots, l_N} \bar{c}_{m_1, \dots, m_N} \left\langle \Omega, \hat{\phi}(f_{k_1})^{l_1+m_1} \cdots \hat{\phi}(f_{k_N})^{l_N+m_N}\Omega \right\rangle_{\mathcal{F}} = \quad (93)$$

$$\stackrel{\text{Step 5}}{\equiv} \sum_{l_1, \dots, l_N=1}^M \sum_{m_1, \dots, m_N=1}^M c_{l_1, \dots, l_N} \bar{c}_{m_1, \dots, m_N} \int_{\mathbb{R}^N} \alpha_{k_1}^{l_1+m_1} \cdots \alpha_{k_N}^{l_N+m_N} \prod_{j=1}^N d\mu_{k_j} \stackrel{\substack{\text{Leb. integr.} \\ \text{linear}}}{\equiv} \int_{\mathbb{R}^N} |P(\alpha_{k_1}, \dots, \alpha_{k_N})|^2 \prod_{j=1}^N d\mu_{k_j}$$

$$\stackrel{\substack{\text{Leb. integr.} \\ \text{linear}}}{\equiv} \int_{\mathbb{R}^\infty} |P(\alpha_{k_1}, \dots, \alpha_{k_N})|^2 d\mu = \int_{\mathbb{R}^\infty} |G^{P;k_1, \dots, k_N}|^2 d\mu = \left\| G^{P;k_1, \dots, k_N} \right\|_{L^2(\mathbb{R}^\infty, d\mu)}^2 =$$

⁴²One could continue until the equality $\star = \prod_{k=1}^N \int_{\mathbb{R}} P_k(\alpha_k) d\mu_k = \prod_{k=1}^N \left\langle 1, P_k(\alpha_k)1 \right\rangle_{L^2(\mathbb{R}^N, d\mu_j)} = \prod_{k=1}^N \left\langle \Omega, P_k(\hat{\phi}(g_k))\Omega \right\rangle_{\mathcal{F}}$.

$$\stackrel{\text{by def}}{=} \left\| S \left(P \left(\hat{\phi}(f_{k_1}), \dots, \hat{\phi}(f_{k_N}) \right) \Omega \right) \right\|_{L^2(\mathbb{R}^\infty, d\mu)}^2.$$

Hence, the operator norm of S is 1.

- **Step 9:** Then, by 3.28 in (Teufel, 2021) (some Hahn-Banach extension theorem), since S is a linear isomorphism between two dense subspaces with operator norm 1, there exists a unique extension operator with the same operator norm, which we call again S , between the closure spaces. That is, there exists a unique unitary $S : \mathcal{F}_+(\mathcal{H}) \rightarrow L^2(\mathbb{R}^\infty, \mathbb{C}.d\mu)$ extending the map in (92). By construction we have that $S\hat{\phi}(f_k)S^{-1} = \alpha_k$ and $(S\Omega)(\alpha) \equiv 1$. o.ε.δ.

Therefore, what we called single boson states in Fock space, i.e., those spanned by linear combinations of $\hat{\phi}(f_j)\Omega$ are mapped to linear combinations of polynomials $\psi(\alpha) = \alpha_j\Psi_0$ (note that $\Psi_0(\alpha) \equiv 1$, but we keep a ghost of it to remember that *heuristically* it corresponds to the Gaussian $e^{-\sum_{j=1}^\infty |\alpha_j|^2}$ that we pushed to the measure). As such, a single boson state in the functional representation is a first order polynomial of arbitrary variables. Likewise, two boson states $\sum_{j,k} c_{j,k}\hat{\phi}(f_k)\hat{\phi}(f_j)\Omega$ are identified with second order polynomials of possibly infinitely many variables $\sum_{j,k} c_{j,k}\alpha_j\alpha_k\Psi_0$. And so on.

To shed a light on what we found and link it with all the previous, denote by H_n^k the n -th Hermite polynomial along the axis α_k , i.e., $H_n^k(\alpha) := H_n(\alpha_k)$. Then, the “creation of a boson” from the vacuum Ω to a state f_k of the ONB, which is given by the Fock creation operator as $a_+^*(f_k)\Omega$, is mapped into the promotion of the “ground state” wavefunctional Ψ_0 to the “excited” state $\Psi_{k\text{-th dof 1 excit, the rest 0}} := H_1^k\Psi_0$. Likewise, the creation of say, 3 bosons in state f_k and 5 in state f_j , namely $a_+(f_k)^3 a_+(f_j)^5 \Omega$ is mapped to $\Psi_{k\text{-th dof 3 excit, j-th 5 excit, rest 0}} := H_3^k H_5^j \Psi_0$. Hence, as promised, we are playing the same game as in §D.3.3, but now with infinitely many modes that can take excitations on them.

Finally, the set of possible combinations of bosons in states of the ONB $\mathcal{B} := \{f_j\}_{j \in \mathbb{N}}$ that we employed to build the unitary S , namely,

$$\left\{ \frac{1}{\sqrt{n_1! \dots n_N!}} a_+^*(f_{j_1})^{n_1} \dots a_+^*(f_{j_N})^{n_N} \Omega \mid N, n_k \in \mathbb{N} \text{ and } f_{j_k} \in B \text{ and } j_1 < \dots < j_N \right\}$$

is an ONB of the Fock space $\mathcal{F}_+(\mathcal{H})$ (by Proposition 5). These are mapped unitarily to products of the Hermite polynomials

$$\{H_{n_1}^{j_1} \dots H_{n_N}^{j_N} \mid N, n_k \in \mathbb{N} \text{ and } j_1 < \dots < j_N\},$$

which are still an ONB, now of $L^2(\mathbb{R}^\infty, \mathbb{C}, d\mu)$.

With all, the bosonic creation and annihilation operators of Fock space, still correspond to ladder operators in infinite dimensions.

D.4.1. A Rigorous Pilot-Wave Theory in \mathbb{R}^∞ ?

To be continued in the Master’s thesis next semester...

...

E. A MEDITATION IN THE SUMMIT: “THERE WERE HIGHER PEAKS ABOVE THE CEILING...”

Now that we finally arrived to the “ceiling” of the cut-off expansion sequence (for the free theory), taking advantage of the “sights from here”, let us close the manuscript with some observations.

- Certainly, there are still many open questions regarding to what extent the wavefunctional QFT we found can be understood as a QFT over the space of mode expansions. Everything points to that, but as we found earlier, \mathbb{R}^∞ has in it more than expansion coefficients (more points than those in ℓ^2) and as the Wiener measures showed in the Feynman-Kac construction, it might very well be the case that our preferred subspace, ℓ^2 , has measure 0 altogether. If so, cleverer ideas will need to be thought, perhaps in the direction of distributions instead of fields. Perhaps, just like in the case of the Feynman-Kac formula, we might need to face that only kinky fields, say, nowhere differentiable ones, are the typical ones. However, as [Dürr and Lazarovici \(2020\)](#) posed it, this feels like a *déjà vu*: recall the time when we decided to leave rational numbers to go to real numbers, just because maths got more “beautiful” in doing so. Not that we really needed them for making experimental predictions, but that way we would avoid talking about infinite cut-off towers of fractions (familiar...). But then we found ourselves in a hurry to define a new measure. The counting measure so perfectly meaningful to intuition was meaningless in the continuum and needed to be substituted by a new measure. We sadly said goodbye to the nice properties of the counting measure and embraced the awkward new measure: the Lebesgue measure. But the most dramatic part of it was still about to come (just like it happened to us...): our beautiful rational numbers suddenly were lost in measure and typicality. All of them together were made to weigh zero, and almost every number according to the new strange measure were numbers with crazy random sequences of decimal numbers. Today we know that in fact that almost all of them are not even computable! But hey, the measure of open balls was made finite in turn (so familiar indeed!).
- Regarding whether we really achieved the ceiling of the cut-off tower, for free theories it seems so. However, as it can be seen explicitly for the case of fermion QFTs ([Deckert et al., 2010](#)), and more in general by the so-called “Haag’s theorem”, although QFTs with no cut-offs can nicely be made rigorous for free species of particles using Fock space techniques, there is no way to make a rigorous mathematical description of interactions using the free particle Fock QFTs (unless we only want to talk about asymptotically free states and scattering matrices). And therefore, by unitary equivalence, neither can be the case if we employ the wavefunctional space we just built. In fact, as of today, every interacting QFT in 4D spacetime (so the whole Standard Model) is “pre-mathematical” unless we accept the cut-off towers as satisfactory theories.

Still, it is known ([Deckert et al., 2010](#)) that part of the problem with interactions and Fock spaces is that interactions force us to consider different (even inequivalent) Fock spaces at each time, tightly related to the fact that the vacuum (say the Dirac sea) need to be re-defined (infinitely many virtual particles emerge from the vacuum). Now, we found that in the discretized versions of the wavefunctional theory we have built (where there was still a Lebesgue measure), there was nowhere a preferred state to be called “vacuum”. Only the identification with Fock space gave so. It is also true that if we had a Hamiltonian with a ground-state then there would be a distinguished state, but there would be no problem for such a ground state to change in time, unlike in Fock space, where the vacuum is given by a fixed vector. Now, in the “no cut-off” limit, we designed the Gaussian measure so that the state equal to 1 everywhere was made to represent the Fock vacuum. Hence, we forced the resulting limit wavefunctional space to have a distinguished state, which was moreover encoded in the measure. But that was arguably only with the purpose of giving it a Fock space look in the search of a unitary equivalence. There are several physicists, like [Sebens \(2022\)](#) or [Jackiw \(1995\)](#) that hope the property of the cut-offs, where different “vacuums” can be described in the very same Hilbert space, will soon be translated into an infinite dimensional space. That is, they hope that there exists a way to talk about different vacuums and perhaps inequivalent Fock spaces as parts of the same wavefunctional space. Indeed, in their heuristic formulas, they already assume so. This would be an amazing step to solve the problem of interactions in QFT and hence, worth to be further studied.

- Another advantage of the wavefunctional approach is that actually the Standard Model is built on the Lagrangian QFT style of heuristic formulas (the interaction Hamiltonians etc.), which are simpler to translate to wavefunctional language than to Fock space. However, it must be conceded that most verifications of experimental results, which are done using scattering matrices, cross sections and asymptotically free QFTs, are commonly done in (more or less heuristic) Fock space pictures. And the problem is that it is hard to

see where the probabilistic interpretation for the field configuration-space, that we fought so hard to make rigorous, fits with the probabilistic interpretation of Fock space. The point is the following: maybe bosons are merely virtual particles, right, excitations of the underlying ontological fields, but if so, we need a mechanism to go from probabilities to find n particles in a configuration (x_1, \dots, x_n) as given by a Fock wavefunction $|\psi_t^{(n)}(x_1, x_2, \dots, x_n)|^2 d^n x$ to the probabilities to find the ontological fields in a configuration with coefficients $(\alpha_1, \alpha_2, \dots)$, as given by the wavefunctional $|\Psi_t(\alpha_1, \alpha_2, \dots)|^2 d\mu$. A heuristic discussion in this direction can be found in (Struyve, 2010).

- It is apparent at this point that QFT is a genuine source of problems, which appeal all disciplines, from experiments to philosophy. Of course, one could be happy with cut-off theories and “renormalization” and believe the tower of approximate QFTs is all there is. Certainly, the trick will continue to work as we keep making the predictive algorithms, the number of parameters and ad-hoc tweaks increase. But at some point we will need to consider seriously (and beware, because the time is about to come) which is the difference between fitting a model to the experimental data (hence, being able to predict phenomenology) just like any neural network is capable of doing (with enough weights and tweaks) and doing physics. We claim that physics is not **fitting and predicting** “at all costs”. Instead, it is the development of a concise mathematical and mechanistic narrative to **understand** the edge between phenomenology and whatever lays beyond it. As a necessary by-product it indeed allows the prediction of phenomenology. But, we claim that obviously, this is not its only quiddity.

Let us close with an adapted version of a fine “bar” by Dürr and Lazarovici (2020) in this direction: “The infinities which appear abundantly in” QFT cannot be “merely mathematical problems that we can try to solve by new techniques” (as both particle physicists and mathematicians in AQFT would like to believe). “They are clear signs of fundamental physical problems which will lead to a distortion of the notion of physical theory if we keep trying to push them to one side: non-convergent perturbative expansions and unspeakable limits of renormalization procedures become the substitute for perhaps two or three fundamental equations” and a clear picture of the world. To which we add: let us not *renormalize* what we mean by *physical theory*!

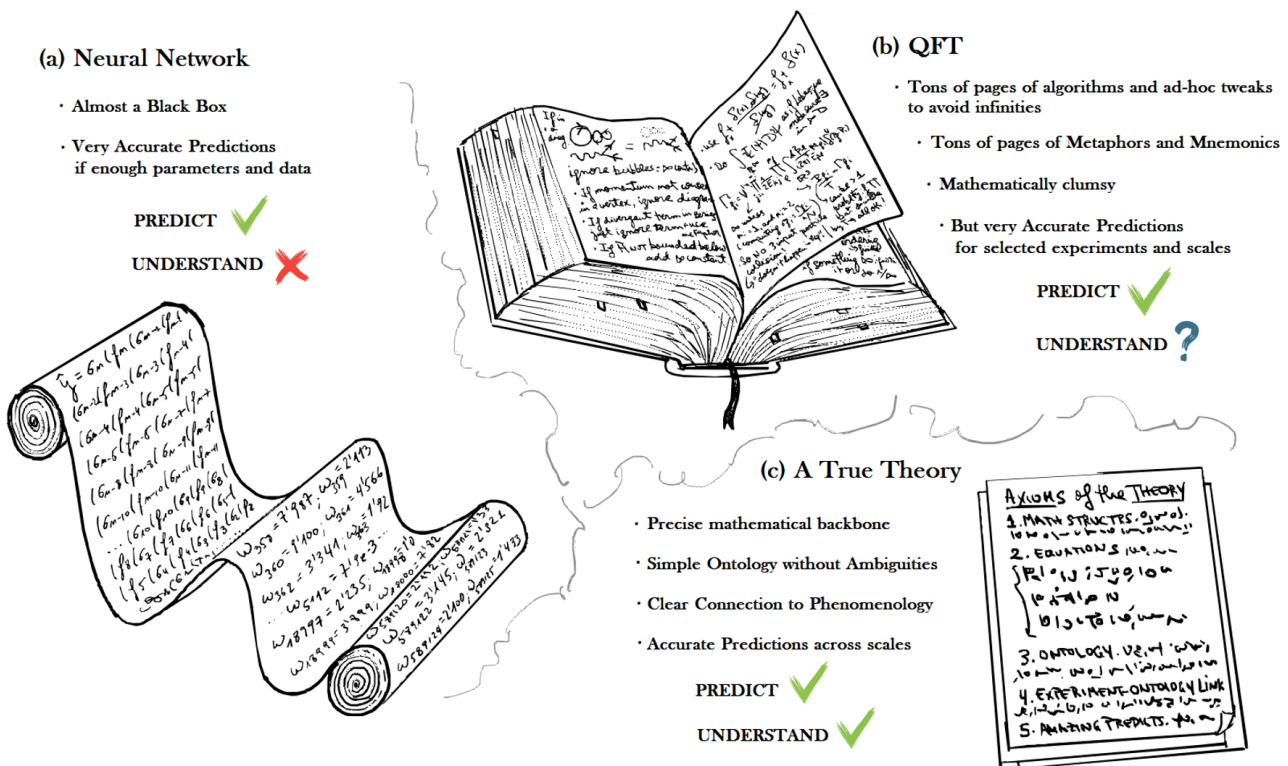


Figure 3: A satirical sketch by the author describing the deformation (or perhaps “renormalization”) of the notion of “physical theory”. As of 2024, we are still in (b) so perhaps we are still in time to go back to (c).

F. BIBLIOGRAPHY

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G. APPENDICES

Appendix to §B.1: Calculus in Infinite Dimensional Vector Spaces

This section drinks mainly from chapter 2 in the highly recommended book by [Abraham et al. \(2012\)](#). Recall that a normed vector space (NVS) is a vector space equipped with a norm, which induces a distance function and hence, a ball topology, i.e., it is also a metric space. We call it Banach if the topology is Cauchy complete. In particular any inner product vector space (IPVS), which is a vector space with an inner product, has a natural norm and hence is a NVS. We call it Hilbert if Cauchy complete. We say a NVS is finite dimensional (e.g., $(\mathbb{R}^n, \|\cdot\|_2)$ or $(\mathbb{C}, |\cdot|)$) if there exists a set of finitely many linearly independent elements spanning the whole space (any element is writeable as a finite linear combination of them), else it is infinite dimensional. Examples of infinite dimensional NVS include $(L^2(\mathbb{R}^n), \|\cdot\|_{L^2})$ or the space of continuous linear maps between two NVS E, F in the operator norm, denoted by $(\mathcal{L}(E, F), \|\cdot\|_{op})$. Hence, a (wave)functional $f : L^2(\mathbb{R}^n) \rightarrow \mathbb{C}$, element of the functional space $L^2(L^2(\mathbb{R}^n), \mathbb{C}, d\mu)$, is nothing but a map between two NVS. Think of such functional spaces as examples for the following definitions.

Definition S1. Let E, F be NVS, $U \subseteq E$ open and $f : U \rightarrow F$ a mapping. We say that f is **differentiable** at $u_0 \in U$ when there exists a continuous linear map $Df|_{u_0} \in \mathcal{L}(E, F)$ such that

$$\lim_{u \rightarrow u_0} \frac{f(u) - f(u_0) - Df|_{u_0}(u - u_0)}{\|u - u_0\|} = 0. \quad (94)$$

If it exists, there is a unique such linear map (hence the notation $Df|_{u_0}$) and we call it the derivative or (total) **differential** of f at u_0 . It is the well-known definition from first year calculus, but now also working for infinite dimensional vector spaces (for whom we call it *Fréchet derivative*).

Definition S2. In the notation, of Def. S1.,

- If f is differentiable at each $u_0 \in U$ then $Df : U \rightarrow \mathcal{L}(E, F)$ is called the (total Fréchet) derivative or differential of f . If it is a continuous map, we say f is **continuously differentiable** and denote the set of continuously differentiable maps by $\mathcal{C}^1(E, F)$.
- We can then define (if it exists) the r -th differential $D^r f := D(D^{r-1} f) : U \subset E \rightarrow \mathcal{L}(E, \mathcal{L}(E^{r-1}, F)) \simeq \mathcal{L}(E^r, F)$, such that $D^r f|_{u_0}$ is an r -linear continuous map for each $u_0 \in U$. If $D^r f$ is continuous, we say f is **r times continuously differentiable** or of **class $\mathcal{C}^r(E, F)$** .

Definition S3. In the notation of Def. S1, we say that f has a **directional derivative in direction $v \in E$ at the point $u \in U$** when

$$\frac{d}{dt} f(u + tv)|_{t=0} := \lim_{t \rightarrow 0} \frac{f(u + tv) - f(u)}{t} \text{ exists.} \quad (95)$$

Proposition S1. In the notation of Def. S1,

- (i) If f is differentiable at $u \implies$ all directional derivatives exist and

$$\frac{d}{dt} f(u + tv)|_{t=0} = Df|_u v. \quad (96)$$

- (ii) $f \in \mathcal{C}^1(U, F) \iff$ for all $v \in E$, the directional derivative exists at all points $u \in U$ and are continuous functions of u for fixed $v \in E$.
- (iii) If $f \in \mathcal{C}^r(U, F)$, for all $v_1, \dots, v_n \in E$ and $u \in U$

$$D^r f|_u(v_1, \dots, v_n) = \frac{d}{dt_1} \cdots \frac{d}{dt_n} f\left(u + \sum_j t_j v_j\right)|_{t_1=\dots=t_n=0}. \quad (97)$$

Note that many multivariate calculus results still hold almost verbatim in arbitrary NV ([Abraham et al., 2012](#)). Among others:

- The Schwarz cross derivative rule: $f \in \mathcal{C}^r \Rightarrow D^r f|_u$ is a symmetric r -linear map in $\mathcal{L}(E^r, F)$.
- Taylor's expansion theorem: If $U \subseteq E$ is convex and $f \in \mathcal{C}^r(U, F)$, then, for all the $(u, v) \in U \times E$ such that $u + v \in U$, there exists a symmetric $R|_{(u,v)} \in \mathcal{L}(E^r, F)$ such that

$$f(u + h) = f(u) + \frac{Df|_u}{1!}h + \frac{Df^2|_u}{2!}(h, h) + \cdots + \frac{Df^r|_u}{r!} \underbrace{(h, \dots, h)}_{r \text{ times}} + R|_{(u,v)} \underbrace{(h, \dots, h)}_{r \text{ times}} \quad (98)$$

with $h := v - u$. Moreover, $R|_{(u,v)}(h, \dots, h) = o(\|h\|_E^r)$.⁴³

- $f : U \subseteq E \rightarrow \mathbb{R}$ is differentiable in $u \in U$ and has a local extremum in $u \in U \implies Df|_u = 0$.
- Results like the sufficient conditions for local maxima or minima, the inverse function theorem, the implicit function theorem, differentiation under the integral sign etc., also hold almost verbatim.

Then, the *functional derivative* is nothing but the gradient related to the differential of a map by some inner product, or more generally by some pairing (a slightly weaker notion). Let us explain this more in detail.

Definition S4. Let E, F be NVS.

- A **pairing (between elements of E and F)** is a continuous, bilinear functional $\langle \cdot, \cdot \rangle : E \times F \rightarrow \mathbb{R}$. (For example if $E = F = L^2(\mathbb{R}^d, \mathbb{R})$, the inner product of L^2 is a pairing).
- A pairing is called **E -non-degenerate** when $\langle x, y \rangle = 0 \forall y \in F \implies x = 0$. That is, when $x \mapsto \langle x, \cdot \rangle$ is an injective map $E \rightarrow F^*$.

For example if $E = F = L^2(\mathbb{R}^d, \mathbb{R})$, the inner product of L^2 is non-degenerate in both slots by the Riesz representation theorem. But if we restrict the second arguments to only $\mathcal{C}_c^\infty(\mathbb{R}^d)$ (smooth compactly supported functions), still it is L^2 non-degenerate by the du Bois-Reymond lemma of PDE theory: i.e,

$$\text{if } \psi \in L^2(\mathbb{R}^d, \mathbb{R}) \text{ is such that } \int_{\mathbb{R}^d} \psi(x)\eta(x)d^n x = 0 \quad \forall \eta \in \mathcal{C}_c^\infty(\mathbb{R}^n) \implies \psi = 0. \quad (99)$$

We defined the general notion of pairing so we can find functional derivatives using the inner product of some ambient space like L^2 in smaller subspaces like \mathcal{C}_c^∞ . And we could do so in both slots: for instance, we could take the L^2 inner product restricted as before, to $\mathcal{C}_c^\infty(\mathbb{R}^d)$ in the second slot, but to $H^2(\mathbb{R}^3, \mathbb{R}) \subset L^2(\mathbb{R}^3, \mathbb{R})$ in the first slot (the second *Sobolev space*). A priori H^2 has another associated inner product, but by (99) the one of L^2 already gives a non-degenerate pairing with elements of $\mathcal{C}_c^\infty(\mathbb{R})$.

Definition S5. Let E, F be NVS with an E -non-degenerate pairing $\langle \cdot, \cdot \rangle : E \times F \rightarrow \mathbb{R}$ and $f : E \rightarrow \mathbb{R}$ be differentiable at $\varphi \in E$. We define the **functional derivative** of f in φ , denoted by $\frac{\delta f}{\delta \phi}|_\varphi$ to be, if it exists, the unique element of E such that

$$Df|_\varphi \eta = \left\langle \frac{\delta f}{\delta \phi}|_\varphi, \eta \right\rangle \quad \forall \eta \in F \quad \text{i.e.,} \quad Df|_\varphi = \left\langle \frac{\delta f}{\delta \phi}|_\varphi, \cdot \right\rangle. \quad (100)$$

It is also called the **gradient** of f in φ , denoted $grad_{\langle \cdot, \cdot \rangle} f|_\varphi \equiv \frac{\delta f}{\delta \phi}|_\varphi$. If E is a function space, we denote $\frac{\delta f}{\delta \phi(x)}|_\varphi := \frac{\delta f}{\delta \phi}|_\varphi(x)$.

Hence, the functional derivative is the generalization of the gradient and is a mapping from points in $U \subseteq E$ to vectors in E (unlike the differential which gives operators, or the directional derivative, which gives vectors of F).

Then, if $f \in \mathcal{C}^1$ we can equivalently characterize functional derivatives by directional derivatives (which yields a simpler formula to compute them). Namely,

$$Df|_\varphi \eta = \frac{d}{dt} f(\varphi + t\eta)|_{t=0} = \left\langle \frac{\delta f}{\delta \phi}|_\varphi, \eta \right\rangle \quad \forall \eta \in F. \quad (101)$$

⁴³Given V is a topological space and W a NVS, for $g : V \rightarrow W$ and $\eta : V \rightarrow \mathbb{R}^+$ we write $g = o(\eta)$ (said, “ g is asymptotically smaller than η ”) when, given an $\alpha \in E$ that is clear from the context, $\lim_{v \rightarrow \alpha} \frac{\|g(v)\|_W}{\eta(v)} = 0$. In the text we have $\alpha = \bar{0} \in E$.

For instance, if $S : L^2(\mathbb{R}^3, \mathbb{R}) \rightarrow \mathbb{R}$ is in \mathcal{C}^1 (and say, represents some action function over fields), then $\frac{\delta S}{\delta \phi}|_\varphi$ is defined by the condition:

$$\frac{d}{dt}S(\varphi + t\eta)|_{t=0} = \int_{\mathbb{R}^n} \frac{\delta S}{\delta \phi(x)}|_\varphi \eta(x) d^n x. \quad (102)$$

In order to find $\frac{\delta S}{\delta \phi}|_\varphi$ we would then develop the left hand side explicitly until we get an expression like the right hand side. By uniqueness of directional derivative we would have found it.

Now, we say that S has a **critical point** at $\varphi \in L^2(\mathbb{R}^3, \mathbb{R})$, when $Df|_\varphi = 0$ (which is a necessary condition for local extremum). If the gradient exists for some non-degenerate $\langle \cdot, \cdot \rangle$, a critical point will equivalently be characterized by $\frac{\delta f}{\delta \phi}|_\varphi = 0$. This equivalence is given by noting that

$$DS|_\varphi = 0 \iff \frac{d}{dt}S(\varphi + t\eta)|_{t=0} = 0 \forall \eta \in \mathcal{C}_c^\infty \iff \frac{\delta S}{\delta \phi}|_\varphi = 0. \quad (103)$$

Now, before proving the field theoretic Euler-Lagrange equations as an example use-case of all this machinery, we need a last ingredient: is there a notion of ‘‘partial derivative’’ with respect to only some arguments?

Definition S6. Let $E_1, \dots, E_n, F_1, \dots, F_n$ be NVS, $f : E_1 \times \dots \times E_n \rightarrow \mathbb{R}$ be in \mathcal{C}^1 and $\langle \cdot, \cdot \rangle_j : E_j \times F_j \rightarrow \mathbb{R}$ be E_j -non-degenerate pairings. We define the **(partial) functional derivative with respect to ϕ_j at the point $\vec{\varphi} := (\varphi_1, \dots, \varphi_n) \in \prod_{j=1}^n E_n$** as the unique element of E_j (if it exists), denoted by $\frac{\delta f}{\delta \phi_j}|_{\vec{\varphi}}$ such that

$$D_j f|_{\vec{\varphi}} \cdot \eta_j := Df|_{\vec{\varphi}}(0, \dots, 0, \eta_j, 0, \dots, 0) = \frac{d}{d\varepsilon}|_{\varepsilon=0} f(\varphi_1, \dots, \varphi_j + \varepsilon \eta_j, \dots, \varphi_n) = \left\langle \frac{\delta f}{\delta \phi_j}|_{\vec{\varphi}}, \eta_j \right\rangle_j. \quad (104)$$

Example: The Field Euler-Lagrange Equations. Let there be some open set $\Omega \subset \mathbb{R}^n$ with compact closure (\mathbb{R}^n can include a time direction, so it could be a space-time) and let $M := H^2(\Omega, \mathbb{R})$ denote the NVS of fields over (the piece of space-time) $\Omega \subset \mathbb{R}^n$ that are square integrable and twice (weakly) differentiable with square integrable (weak) derivatives. Now, let there be a function $S : M^m \rightarrow \mathbb{R}$, which we call the **action**, defined by

$$S(\varphi_1, \dots, \varphi_n) := \int_{\Omega} \mathcal{L} \left(x, \varphi_j(x), \frac{\partial \varphi_j}{\partial x_k}(x) \right) d^n x \quad (105)$$

for some smooth function $\mathcal{L} : \Omega \times \mathbb{R}^m \times \mathbb{R}^{n \times m} \rightarrow \mathbb{R}$, of arguments $\mathcal{L}(x, \phi_j, \frac{\partial \phi_j}{\partial x_r})$, called the **Lagrangian density**. According to the postulate known as **action principle**, the actual fields $\vec{\varphi}$ we observe in the world are critical points of the action S . But as we saw above, $\vec{\varphi}$ is a critical point of S if and only if for some pairing for which the functional derivative exists, $\frac{\delta S}{\delta \phi_k}|_{\vec{\varphi}} = 0$ for all $k \in \{1, \dots, m\}$. Hence, if we find the functional derivatives, we will obtain an explicit equation governing the critical points of S .

A simple way to find a functional derivative with respect to the k -th argument in the case of an action-like function (given by the integration of some density), is to take the directional derivative in an arbitrary direction $\eta \in \mathcal{C}_c^\infty(\Omega)$ along the k -th argument and to manipulate the expression until we get the L^2 pairing between the η and some other vector ψ . By the uniqueness of functional derivative, the ψ must be the searched functional derivative. As a showcase, let $\eta_k \in \mathcal{C}_c^\infty(\Omega, \mathbb{R})$ and fix a point in the configuration space(-time) $\vec{\varphi} := (\varphi_1, \dots, \varphi_m) \in M^m$, then,

$$\begin{aligned} D_k S|_{\vec{\varphi}} \eta_k &:= \frac{d}{d\varepsilon}|_{\varepsilon=0} S(\varphi_k + \varepsilon \eta_k; \{\varphi_j\}_{j \neq k}) = & (106) \\ &= \frac{d}{d\varepsilon}|_{\varepsilon=0} \int_{\Omega} \mathcal{L} \left(x, \varphi_k(x) + \varepsilon \eta_k(x), \frac{\partial \varphi_k(x)}{\partial x_r} + \varepsilon \frac{\partial \eta_k(x)}{\partial x_r}; \left\{ \varphi_j(x), \frac{\partial \varphi_j(x)}{\partial x_r} \right\}_{j \neq k} \right) d^n x = \\ &= \underbrace{\int_{\Omega} \left[\frac{\partial \mathcal{L}}{\partial \phi_k}|_{\vec{\varphi}}(x) - \sum_{r=1}^n \frac{\partial}{\partial x_r} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \phi_k}{\partial x_r} \right)}|_{\vec{\varphi}}(x) \right) \right]}_{\substack{\text{different. under integral +} \\ \text{chain \& Leibniz rules +} \\ \eta_k(\partial\Omega)=0 \text{ bec. comp. supp}}} \eta_k(x) d^n x = \\ &= \left\langle \frac{\partial \mathcal{L}}{\partial \phi_k}|_{\vec{\varphi}} - \sum_{r=1}^n \frac{\partial}{\partial x_r} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \phi_k}{\partial x_r} \right)}|_{\vec{\varphi}} \right), \eta_k \right\rangle_{L^2(\Omega, \mathbb{R})} =: \left\langle \frac{\delta S}{\delta \phi_k}, \eta_k \right\rangle_{L^2}. \end{aligned}$$

Hence, $\frac{\delta S}{\delta \phi_k} |_{\vec{\varphi}} = 0$ for all $k \in \{1, \dots, m\}$ if and only if

$$\frac{\partial \mathcal{L}}{\partial \phi_k} |_{\vec{\varphi}} - \sum_{r=1}^n \frac{\partial}{\partial x_r} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \phi_k}{\partial x_r} \right)} |_{\vec{\varphi}} \right) = 0 \quad \text{for all } k \in \{1, \dots, m\}. \quad (107)$$

This system of PDEs is called **the Euler-Lagrange equations**.

Appendix to §B.2: Not all PWTs explain the collapse and “measurement” probabilities!

It is not at all clear that a non-position configuration pilot-wave theory,⁴⁴ in the context of a quantum “measurement”, can predict without postulating it, the collapse of the wavefunction or the probabilities associated to its outcomes. And like it or not, both are experimentally verified facts. Let us find why this is not clear by reviewing the key aspects of why Bohmian mechanics does predict them. (Spoiler: because its configuration space is over **position-space!**).

• First Act: An Interpretation Agnostic Setting of an ideal Quantum Measurement

Let Q be the configuration-space of the laboratory (or if preferred, the whole Universe), whatever it is: position configuration or others and either particle or field configuration-space. As an obvious assumption in physics, we can split it as $Q = Q_{object} \times Q_{device}$, with Q_{object} the configuration-space of an *object-system* we want to study and Q_{device} containing the rest of degrees of freedom (in particular, the devices of the laboratory).⁴⁵ We assume the object is sufficiently isolated from the rest of the environment to have “its own unitary SE”, meaning that it has some normalized effective wavefunction $\psi \in L^2(Q_{object})$ (an obvious assumption in any quantum treatment where we want to consider only the object system with no external influence).

Now, assume we are interested on a countable (or finite) set of orthonormal wavefunctions $\{\varphi_j\}_{j \in J} \subset L^2(Q_{object})$ (e.g., eigenfunctions of the object system’s Hamiltonian). We can bookkeep them in the self-adjoint operator $\hat{D} := \sum_{j \in J} j \varphi_j \langle \varphi_j, \cdot \rangle$ —so that we can recover them up to global phase from the spectral decomposition of \hat{D} . We build a (macroscopic) device (i.e., such that we can distinguish its different dial positions with the “naked eye”). The configuration of the device is described within Q_{device} . We build the device with a dial that can take the values in J and is built so that **if** $\psi \in \{\varphi_j\}_{j \in J}$, after the device is brought to interact with the object, we can unambiguously determine in which of those states was the object prior to the interaction. We call this assumption (\star), telling that by calibration/design of the device:

- (\star) Whenever the object’s ψ is prepared so that it is proportional to φ_j (i.e., it is in the eigenspace of eigenvalue j of the bookkeeping s.a. operator \hat{D}), after the device-object coupling, the dial of the device shows the direction $j \in J$ almost surely (probability 1), i.e., it happens so in (almost) all repetitions of the experiment.

Arguably, this qualifies the device to be called \hat{D} -“measurement” device, but in general it is rather called \hat{D} -“quantum measurement” device (because as we will mention later, the object can also initially be in a superposition of the states in $\psi \in \{\varphi_j\}_{j \in J}$ and then...we will see what happens).

If initially the object has an effective wavefunction because it is sufficiently isolated from the devices, so must its complement (the devices) have by reciprocal isolation from the object. We call it $\phi_R \in L^2(Q_{device})$, which we will assume to be compatible with the “Ready to take measurements”-state of the device.

• Second Act: Interpretation Agnostic Dynamics of an ideal Quantum Measurement

For any preparation $\psi \in L^2(Q_{object})$ (among the eigenstates of \hat{D} or not), the process of “measurement” starts with a joint device-object wavefunction $\phi_R \otimes \psi \in L^2(Q_{device} \times Q_{object})$. As an obvious fact, for the extraction of information from the object there must be an influence of the object on the device, meaning that the device will need to break the isolation of the object system by interacting with it for some time $\tau > 0$.⁴⁶ By

⁴⁴Meaning that the primitive ontology is not a trajectory unambiguously translatable to position-space. An example would be a canonical momentum trajectory.

⁴⁵For example, if we chose $Q = \mathbb{R}^{3N}$, with $N \in \mathbb{N}$ the number of particles in the lab (or the Universe), such that n of those particles belong to the object system, then $Q_{object} = \mathbb{R}^{3n}$. If instead, we assumed Q to be position-space fields, say $Q = H^2(\mathbb{R}^3, \mathbb{R})$, then Q_{object} would be the set of position-space fields localized in the object’s domain $\Omega \subset \mathbb{R}^3$, say $H^2(\Omega, \mathbb{R})$.

⁴⁶As always in the known physics, by mutuality of interactions, an interaction will not only cause the object system to influence the device and its pointer, but this will reciprocally influence/perturb/back-act the object system.

the Schrödinger equation dictating the joint dynamics, the coupling happens through a unitary propagator U^τ acting on $L^2(Q_{device} \times Q_{object})$ that will leave towards the end of the measurement a joint state $U^\tau(\phi_R \otimes \psi)$. This is generally an entangled device-object wavefunction.

In particular, (\star) and the linearity of the Schrödinger equation (thus, linearity of the unitary U^τ) imply

$(\star\star)$ If the object system is prepared in a superposition of the states for which our device was built, say, if $\psi = \sum_{j \in J} c_j \varphi_j$ for $\{c_j\}_j \subset \mathbb{C}$, then, at the end of the coupling the device-object state is described by the superposition $\sum_{j \in J} c_j U^\tau(\phi_R \otimes \varphi_j)$.

But then we have the experimental constraint that:

$(\star\star\star)$ At the end of the “measurement”, we see that the device points unambiguously to one particular $k \in J$ with a probability $|c_k|^2$.

• **Resolution according to Orthodox QM: an unnecessary and unimaginative “Deus ex machina”**

Let us say that in a particular experiment we get $k \in J$. By calibration of the device, every state $U^\tau(\phi_R \otimes \varphi_j)$ with $j \neq k$ is “incompatible” with the observed macroscopic situation and only $U^\tau(\phi_R \otimes \varphi_k)$ is compatible with it. Then, as one of the most shameless and unprofessional “Deus ex machina”-s in the whole history of science, the textbook QM theory **postulates** that there has been a so-called “collapse of the wavefunction” abruptly taking the superposition $\sum_{j \in J} c_j U^\tau(\phi_R \otimes \varphi_j)$ into $U^\tau(\phi_R \otimes \varphi_k)$, and that this so-happens with probability $|c_k|^2$.

• **Resolution according to Bohmian Mechanics: $(\star\star\star)$ is an obvious consequence of (\star) and the SE**

In Bohmian mechanics, where $Q = \mathbb{R}^{3N}$ and $Q_{object} = \mathbb{R}^{3n}$ are the **position**-configuration space, with $N > 10^{23}$ (because we are describing a whole macroscopic device) and $n < 10^{11}$ (for it is a microscopic object system—a human DNA molecule has around 10^{11} particles), the plot is resolved with no need of divine intervention where “God plays some die”. Within Bohmian mechanics (\star) has the precise meaning that:

(\star) By design of the device, the state $U^\tau(\phi_R \otimes \varphi_j)$ is such that almost all possible device-object Bohmian trajectories are compatible with a macroscopic situation in which the dial shows j .

But then, the following chain of observations and arguments follows:

- The dial must be in macroscopically different positions to show us different j .
- The dial’s position is unambiguously defined by the device-object Bohmian trajectory **because the trajectory lives in position-space**.
- Hence, different dial positions happen if and only if macroscopically different Bohmian trajectories.
- Therefore, by \star , almost every device-object Bohmian trajectory supporting the state $U^\tau(\phi_R \otimes \varphi_j)$ must be macroscopically far away in configuration space from those that support $U^\tau(\phi_R \otimes \varphi_k)$ (with $k \neq j$).
- Finally, because the typicality measure of Bohmian trajectories is exactly the magnitude (squared) of the wavefunction, the Bohmian trajectories supporting a state make up its functional support, so that:

$(\star\star)$ The supports of $U^\tau(\phi_R \otimes \varphi_j)$ and $U^\tau(\phi_R \otimes \varphi_k)$ must be macroscopically disjoint in Q , i.e., a macroscopic distance (several orders of magnitude more than the characteristic dispersion-scales of φ_j) must separate their supports. We say they are macroscopically disjoint wavepackets.

So, we proved that in Bohmian mechanics, $(\star) \Rightarrow (\star\star)$.

On the other hand, $(\star\star)$ and $(\star\star)$ together mean that at the end of the measurement, the joint state $\psi^\tau := \sum_{j \in J} c_j U^\tau(\phi_R \otimes \varphi_j)$ is the sum of macroscopically disjoint wavepackets. By the dispersive time evolution of the Schrödinger equation, macroscopically disjoint wavepackets made of a macroscopic amount of particles remain macroscopically disjoint for almost always (this loss of capability to self-interfere again, for the macro-disjoint portions of a many-particle wavefunction, is called “*decoherence*”). As such, after the “measurement”, the packets $c_j U^\tau \phi_R \otimes \varphi_j$ evolve independently of each other in the joint configuration-space. But now, in Bohmian mechanics there is a determinate position for all the particles at all times and they follow the typicality measure of the magnitude squared of the wavefunction, so the device-object Bohmian trajectory will end up in one of the supports of the wavepackets, say, in the k -th one. This explains why we see a determinate dial at all times that ends up pointing in some k -th direction (in agreement with the experimental constraint $(\star\star\star)$). Moreover, since the different $c_j U^\tau(\phi_R \otimes \varphi_j)$ in the superposition evolve independently of each other, effectively

the device-object wavefunction is $c_k U^\tau \phi_R \otimes \varphi_k$ (explaining why the collapse postulate of Orthodox QM is indeed acceptable to be assumed effectively).

Finally, denote by $y \in \mathbb{R}^{3m}$ the configuration of the device (including the dial) and $x \in \mathbb{R}^{3n}$ the configuration of the object. Define $B_k \subset \mathbb{R}^{3m}$ as the set of configurations of the device corresponding to a dial pointing in the k -th direction. By $(\star\star)$,

$$\text{supp } U^\tau(\phi_R \otimes \varphi_k) \cap B_k = \begin{cases} \emptyset & \text{if } j \neq k \\ \text{supp } U^\tau(\phi_R \otimes \varphi_k) & \text{if } j = k \end{cases}. \quad (108)$$

Hence, the probability that at the end of the ‘‘measurement’’ the dial points in the k -th direction, according to Bohmian mechanics, is:

$$\begin{aligned} \text{Prob}(y \in B_k) &= \int_{B_k} \int_{\mathbb{R}^n} |\psi^\tau(x, y)|^2 d^n x d^m y = \int_{B_k} \int_{\mathbb{R}^n} \sum_{j \in J} |c_j U^\tau \phi_R \otimes \varphi_j|^2 d^n x d^m y = \\ &\stackrel{(108)}{=} \int_{\text{supp } U^\tau(\phi_R \otimes \varphi_k)} \int_{\mathbb{R}^n} |c_k|^2 |U^\tau \phi_R \otimes \varphi_k|^2 d^n x d^m y \stackrel{U^\tau \text{ is unitary}}{=} |c_k|^2. \end{aligned} \quad (109)$$

Exactly as found experimentally in $(\star\star\star)$.

• Necessary Conditions for a Pilot Wave Theory to use the Bohmian Resolution

Proposition S2. Given a pilot wave theory, the following is a necessary condition to be able to use the same proof as Bohmian mechanics to derive from the postulates of the pilot-wave theory, the collapse of the wavefunction and the probabilistic outcomes of ideal quantum measurements (by just substituting ‘‘Bohmian trajectory’’ by ‘‘primitive ontology’s trajectory’’):

The macroscopic position of a dial is unambiguously defined by the trajectory of the device-object primitive ontology.

The key reason why this holds is that in physics, the image of the macroscopic world is based on positions,⁴⁷ so that different macroscopic situations can only be unambiguously given by position-configurations of it. In very simple terms, knowing for instance the classical canonical momentum trajectory of a set of particles does not tell us how far apart their associated position trajectories could be, hence, we could not tell if the dial says this or that if we only a momentum trajectory for it. The problem is that we could attribute more than one position trajectory to the same momentum trajectory. With this in mind, the obvious quality that tells us if a dial is pointing in this or that direction with least doubt, is the position (and subordinate coarse-grainings of it).

Examples that fall in the necessary condition are the pilot wave theories where the primitive ontology has an immediate position representation. For example, if we take fields over position-space $\rho(x, y, z)$ to describe the density of matter in physical space, by using fields that show higher concentrations wherever our macroscopic image of the world goes, we will achieve the above proposition.⁴⁸ Likewise, if we have a field configuration-space of mode expansion coefficients $(\alpha_1, \alpha_2, \dots)$ that represent a density field over position-space via some $\rho(x, y, z) = \sum_j \alpha_j u_j(x, y, z)$, there is a 1 to 1 related position representation that we can use for the Bohmian-like explanation of collapse. Yet another example where there are fields that do not represent densities of matter but still the Proposition S2 holds is the QED we gave in the introduction. There, the device dial would be assumed to be made of electrons and that would be enough. A non-example is a cousin pilot wave theory to Bohmian mechanics where we describe wavefunctions in Fourier space and the Bohmian trajectory is a ‘‘momentum space’’ trajectory. Similar examples could be built for fields over non-position configuration-spaces.

Let us generalize this a bit. The condition of Proposition S2 was necessary to prove $(\star) \Rightarrow (\star\star)$, which was the crucial thing allowing the rest of the derivation of the collapse and its probabilities. But if we find a pilot wave theory that explains $(\star) \Rightarrow (\star\star)$ in an alternative way, we will have also solved the problem.

⁴⁷For a very good reason: whether something occupies a spatial extension or not is the perceptual/phenomenological quality that most humans will agree on. We could take colour, sound, smell etc., but none will be as simply agreeable.

⁴⁸There are some claims that under the Gaussian measures we will develop in the rigorification of the wavefunctional spaces, the set of all wavefunctionals compactly supported in discrete islands (hence giving clearly localized objects) has measure zero. But of course, the fields do not necessarily need to go down to zero nor be differentiable to show well-defined macro-situations. If the difference between the regions of space where they are very dense and those in which they are not is sufficiently large, we will not be able to distinguish the result.

Proposition S3. Given a pilot wave theory, $(\star) \Rightarrow (\star\star)$ is a necessary condition to use, substituting “Bohmian trajectory” by “primitive ontology’s trajectory”, the same proof as Bohmian mechanics (the part after we proved that implication) to derive from the postulates of the pilot-wave theory, the collapse of the wavefunction and the probabilistic outcomes of ideal quantum measurements. ♦

Then, a (stronger) condition to check that indeed, a pilot wave theory is able to explain the collapse and its probabilities would be that:

For any pair of different macroscopic situations A and B , there exist wavefunction(s) $\Psi_A, \Psi_B \in L^2(Q, \mathbb{C}, d\mu)$ in the global (device-object) configuration space, such that:

- *Almost every ontological trajectory piloted by Ψ_A represents the same macroscopic situation A (likewise for B).*
- *Ψ_A and Ψ_B have disjoint support in the configuration-space of primitive ontology Q .*

In our context of QFTs, this would imply to check whether “whenever a wavefunctional yields the same macro-situation for a.e. ontological field, this wavefunctional is disjoint from others yielding other macro-situations for a.e. ontological field.” To check this stronger condition is more conventional in the literature —although it is harder— (see some examples for various QFTs in [Struyve \(2010\)](#)).

Appendix to §D.2: How the Trick is Employed in the Feynman-Kac Formula

In §D.2.1 we suggested a trick that one can employ in the search for a measure that gives a rigorization of Feynman’s path integral over path-space (which is infinite dimensional) for the heat equation. Let us find which is the problem we want to solve with it and how the trick could be employed. What follows drinks from [Reed and Simon \(1975\)](#)’s section⁴⁹ on the topic.

(α) Feynman’s Heuristic

Consider a single Newtonian particle following a trajectory $\tau \mapsto \vec{X}(\tau) \in \mathbb{R}^3$ for $\tau \geq 0$. Given a potential energy field $V : \mathbb{R}^3 \rightarrow \mathbb{R}$ (with suitable regularity), we define the usual action functional

$$S(\vec{X}; t) := \int_0^t \left(\frac{1}{2} m \left| \frac{d\vec{X}(\tau)}{d\tau} \right|^2 - V(\vec{X}(\tau)) \right) d\tau. \quad (110)$$

(We will consider the mass m to be 1/2 hereafter.)

Now, if we wanted to compute such an integral numerically, we would first discretize $[0, t]$, considering a grid of $n \in \mathbb{N}$ intervals in time of equal length $\Delta t := t/n$ and equispaced nodes $t_j := j\Delta t$ for $j \in \{0, \dots, n\}$. Next, we would define the discretization of the trajectory as $x_j := \vec{X}(t_j)$ (see Figure 4.(a)). Employing a trapezoid rule for the integral⁵⁰ and an Euler finite difference⁵¹ for the velocity of the trajectory, we would discretize S as

$$S_n(x_0, \dots, x_n, t) := \sum_{j=0}^{n-1} \frac{t}{n} \left[\frac{1}{4} \left(\frac{|x_{j+1} - x_j|}{t/n} \right)^2 - V(x_j) \right]. \quad (111)$$

Then, by numerical calculus (for differentiable/nice enough \vec{X}), $\lim_{n \rightarrow +\infty} S_n(x_0, \dots, x_n, t) = S(\vec{X}; t)$. Hereafter, we will denote by x the **end-point** of the trajectories, namely, the node $x_n = \vec{X}(t)$.

In quantum mechanics, one can find that if we define $H_0 := -\Delta$ as the single free particle Hamiltonian in $L^2(\mathbb{R}^3)$ (of domain $D(H_0)$),

Theorem S1. For $V \in L^2(\mathbb{R}^3) + L^\infty(\mathbb{R}^3)$, and any initial wavefunction $\psi_0 \in L^2(\mathbb{R}^3)$, the time evolved wavefunction $\psi_t := (e^{-it(H_0+V)}\psi_0)$ is also given by

$$\psi_t(x) = \lim_{n \rightarrow +\infty} \left(\frac{n}{4\pi it} \right)^{\frac{3n}{2}} \int_{\mathbb{R}^3} \dots \int_{\mathbb{R}^3} e^{iS_n(x_0, \dots, x_{n-1}, x, t)} \psi_0(x_0) dx_0 \dots dx_{n-1} \quad (\text{in } L^2 \text{ sense}). \quad (112)$$

⁴⁹Their derivations can be quite confusing for those aiming a rigorification of the physics literature path integral, because they consider that it is the initial point of the paths (and not the final point) what is fixed. We have modified their derivations with care to match the most common physical intuitions.

⁵⁰For a regular enough $f : (0, t) \rightarrow \mathbb{R}$, to leading n order and large $n \in \mathbb{N}$, $\int_0^t f(\tau) d\tau$ is given by $\sum_{j=0}^{n-1} f(t_j) \Delta t$.

⁵¹For a regular enough $t \mapsto \vec{X}(t)$, to leading Δt order in large $n \in \mathbb{N}$, the velocity $\frac{d}{dt} \vec{X}(t_j)$ is given by $\frac{\vec{X}(t_{j+1}) - \vec{X}(t_j)}{\Delta t}$.

◆

Proof. On the one hand, once can find in item 2.19 of (Teufel, 2021) that the free time evolution operator e^{itH_0} is equal to the integral operator:

$$(e^{itH_0}\varphi)(x) = \frac{1}{(4\pi it)^{3/2}} \int_{\mathbb{R}^3} e^{\frac{i|x-y|^2}{4t}} \varphi(y) dy. \quad (113)$$

On the other hand, as found in item 3.124 of (Teufel, 2021) a potential splittable as in the hypothesis makes $H_0 + V$ self-adjoint on $D(H_0)$, so that we can employ the Trotter formula (item 4.36 in (Teufel, 2021)), i.e.,

$$e^{-it(H_0+V)}\psi_0 = \lim_{n \rightarrow +\infty} \left(e^{-i\frac{t}{n}H_0} e^{-i\frac{t}{n}V} \right)^n \psi_0. \quad (114)$$

But in this expression we can evaluate each of the n propagators $e^{-i\frac{t}{n}H_0}$ using (113) in a chained way to get (112) (namely, evaluate (113) first in the left-most free propagator, i.e., take as if $e^{-i\frac{t}{n}V} \left(e^{-i\frac{t}{n}H_0} e^{-i\frac{t}{n}V} \right)^{n-1} \psi_0$ was the φ of (113) and call y instead x_{n-1} ; evaluate the next free propagator, taking as if the φ was now $e^{-i\frac{t}{n}V} \left(e^{-i\frac{t}{n}H_0} e^{-i\frac{t}{n}V} \right)^{n-2} \psi_0$ and using x_{n-2} instead of y . Iterate n times). o.ε.δ.

Let us analyse (112). In the dx_j integral of (112) we are sweeping over all the possible values that a path \vec{X} could take at $\vec{X}(t_j)$. Hence, as we take the limit $n \rightarrow +\infty$, the spatial integrals of the trajectory over the time grid, end up sweeping the values that the trajectory could take at each $\tau \in [0, t)$. The only clear constraint is that the end-point is fixed as $\vec{X}(t) = x$, because we are not integrating over it. This is as if we were integrating the exponential of the action iS over the space of possible paths that end in x (note how this would include all paths, crazily oscillating paths as well). See Figure 4.(a). With all, as the limit exists by (112), and it seems to be interpretable as an integral over the space of paths ending at x , which we denote by Ω_x , the reader must already have on the tip of their tongue that a formula like

$$“\psi_t(x) = \int_{\vec{X} \in \Omega_x} e^{iS(\vec{X}; t)} \psi_0(\vec{X}(0)) d\mu” \quad (115)$$

should exist as the limit of (112). It is said that R. Feynman was the first one having such an insight. But now we know that no matter how much we look for a measure $d\mu$ which is the limit of infinitely many Lebesgue measures “ $d\mu = \lim_{n \rightarrow \infty} dx_0 \cdots dx_{n-1}$ ”, such an object does not exist.

But heuristic (115) would be very nice if true. Among others, it would provide a very fancy way to explain the “classical limit” of QM. If we did not take units where $\hbar = 1$, the propagators would be things like $e^{-i\frac{t}{\hbar}H}$, so that the heuristic would be

$$“\psi_t(x) = \int_{\vec{X} \in \Omega_x} e^{i\frac{S(\vec{X}; t)}{\hbar}} \psi_0(\vec{X}(0)) d\mu.” \quad (116)$$

Then, as \hbar is made smaller and smaller compared to the characteristic values of S (say, because the mass m to which S is proportional, becomes larger and larger as we consider larger systems), the phases would “tend to cancel each other”, except around critical points of S , where S is constant at first order and corresponds to the classical mechanics paths (by virtue of the action principle). So, we would have that in the limit $\hbar \rightarrow 0$, the contribution to dynamics would come exclusively from paths nearer and nearer to the classical path.

Despite physicists use heuristic formula (115) to derive (ill-defined) perturbation series⁵² assuming somebody will make it rigorous somehow, we are still unable to make “Feynman’s dream” come true. What is possible, is to rigorify such a limit for a very closely related expression: for the analogue propagator of the heat equation (with a potential) instead of the Schrödinger equation. The only formal difference between these two equations is that the i multiplying the time derivative in the SE is substituted by a -1 . The point is that this change will be the crucial thing that will allow an immediate usage of the trick in §D.2.1 to find a limit measure for the heat equation’s case.

⁵²They could perhaps be made sense of- by considering discretized versions of such heuristic formulas (e.g., considering (112) to be the discretization of (115)). If so, the heuristic formulas employed in the physics literature would be understandable as symbols representing such discretizations. However, a priori, the derived expressions would be discretization-dependent and it is not clear if they would converge to a common result. Such a situation is analogous to the cut-off techniques mentioned in §C.

(β) A More Promising Path Integral

Let us see how we can use the same proof as that of Theorem 1 to derive an analogous result for the heat equation. One can find in [Teufel \(2021\)](#) (using convolution and Fourier multiplier techniques) that the free heat equation $\frac{d}{dt}\rho_t = \Delta\rho_t$ has a propagator e^{-tH_0} given by a very similar integral operator to that of the Schrödinger equation, so that for $\eta \in L^2(\mathbb{R}^3)$

$$(e^{tH_0}\eta)(x) = \frac{1}{(4\pi t)^{3/2}} \int_{\mathbb{R}^3} e^{-\frac{\|x-y\|^2}{4t}} \eta(y) dy. \quad (117)$$

- Alternatively, one can find this propagator formula using the following insight based on the fact that the heat equation emerges as the time evolution of the probability density of a Brownian particle (as seen in the “Chaos” course). Take a Brownian particle of which we know the $t = 0$ position to be y , so that at $t = 0$ its position is distributed according a Dirac delta measure at y . If we now consider all the possible random walks it could take from there till time t (where each “infinitesimal step” is normally distributed), its position x will follow a normal/Gaussian distribution with density (called “heat-kernel”)

$$K(x; y, t) := \frac{1}{(4\pi t)^{3/2}} e^{-\frac{\|x-y\|^2}{4t}}. \quad (118)$$

Now, if we initially did not know its position y , but instead, it was distributed according to a probability density η , we could predict the probability measure at (x, t) to be the “sum” of the densities arriving to that point via random walks that start at each possible y , weighted by the actual probability density there was at y . This would then leave $\eta_t(x) = \int_{\mathbb{R}^3} K(x; y, t)\eta(y)dy$, which is equation (117). See Figure 4.(c) for a graphical representation of this intuition.

We could then take the Trotter formula and plug the integral operator (117) in it to get a formula akin to the one in Theorem 1. For the Trotter formula we need to assume V is a self-adjoint multiplication operator that is bounded above, with $H := H_0 - V$ essentially self adjoint in $D(H_0) \cap D(V)$. Then, the PDE governing the probability density of a Brownian particle in a potential, which is the heat equation $\frac{d}{dt}\rho_t = -H\rho_t = (\Delta + V)\rho_t$ (i.e., the SE replacing i by -1), has a propagator e^{-tH} (defined using the functional calculus). Moreover, the Trotter formula leaves

$$e^{-tH}\rho_0 = \lim_{n \rightarrow +\infty} \left(e^{-\frac{t}{n}H_0} e^{\frac{t}{n}V} \right)^n \rho_0. \quad (119)$$

Denoting $\rho_t := e^{-tH}\rho_0$, in the last expression we can evaluate (117) for each of the n free propagators (just as done in Theorem 1), to get that rigorously,

$$\begin{aligned} \rho_t(x) &= \lim_{n \rightarrow +\infty} \left(\frac{n}{4\pi t} \right)^{3n/2} \int_{\mathbb{R}^3} \cdots \int_{\mathbb{R}^3} e^{-\sum_{j=0}^{n-1} \frac{|x_{j+1}-x_j|^2}{4t/n}} e^{\sum_{j=0}^{n-1} \frac{t}{n} V(x_j)} \rho_0(x_0) \prod_{j=0}^{n-1} dx_j = \\ &= \lim_{n \rightarrow +\infty} \left(\frac{n}{4\pi t} \right)^{3n/2} \int_{\mathbb{R}^3} \cdots \int_{\mathbb{R}^3} e^{-S_n(x_0, \dots, x_n, t)} \rho_0(x_0) dx_0 \cdots dx_{n-1}. \end{aligned} \quad (120)$$

This is the same provocative formula as (112) but now with a -1 instead of an i . Using the same arguments, we would find on the tip of our tongues the heuristic (115) (now changing i by -1) as the limit expression of (120). Writing S explicitly, such a heuristic would read:

$$“\rho_t(x) = \int_{\vec{X} \in \Omega_x} e^{-\frac{1}{4} \int_0^t \left| \frac{d\vec{X}(\tau)}{d\tau} \right|^2 d\tau} e^{\int_0^t V(\vec{X}(\tau)) d\tau} \rho_0(\vec{X}(0)) \prod_{j=0}^{\infty} dx_j” . \quad (121)$$

And again, there would be no limit measure $\lim_{n \rightarrow +\infty} \prod_{j=0}^{n-1} dx_j$. However, we are now allowed to apply the trick of §D.2: if we rewrite the (rigorous) equation (120) by placing the kinetic energy part next to the Lebesgue measure, we get

$$\rho_t(x) = \lim_{n \rightarrow +\infty} \int_{\mathbb{R}^3} \cdots \int_{\mathbb{R}^3} e^{\sum_{j=0}^{n-1} \frac{t}{n} V(x_j)} \rho_0(x_0) \prod_{j=0}^{n-1} \left(\frac{n}{4\pi t} \right)^{3/2} e^{-\frac{|x_{j+1}-x_j|^2}{4t/n}} dx_j. \quad (122)$$

If we considered instead of dx_j , the measure $d\mu_j := \left(\frac{n}{4\pi t}\right)^{3/2} e^{-\frac{|x_{j+1}-x_j|^2}{4t}} dx_j$, which is a probability measure for fixed x_j (a Gaussian centred in x_j with standard deviation $2t/n$), we know that the limit product measure would now exist! And hence, not the formula (121), but the same formula with the kinetic energy term and the normalization constants absorbed into the measures would indeed have a chance to exist!

The reason why such a trick works here and not in the quantum case can be heuristically understood as follows. In both ideal formulas (115) and (121), differing only in the i changed for -1 , the Lebesgue measure product blows up to infinity (every open set would have infinite measure, as we saw in §D.2). Likewise, for arbitrary paths $\vec{X} \in \Omega_x$ (say, non-differentiable “kinky” paths), the “derivative” $\frac{d\vec{X}(\tau)}{d\tau}$ would blow-up to infinity as well, which would make the integral $\int_0^t \left|\frac{d\vec{X}(\tau)}{d\tau}\right|^2 d\tau$ highly singular. However, the exponential with the minus sign in (121) would make the latter blow-up arbitrarily close to zero, and “lumping” together this zero of the kinetic energy with the blowing up Lebesgue measures, gives a well-defined limit measure. In the quantum case (115) on the other hand, since the integral of the kinetic energy in the exponential does not have a minus sign, but instead an i , it does not decay for singular derivatives and there is no term to compensate the blowing up Lebesgue measure.

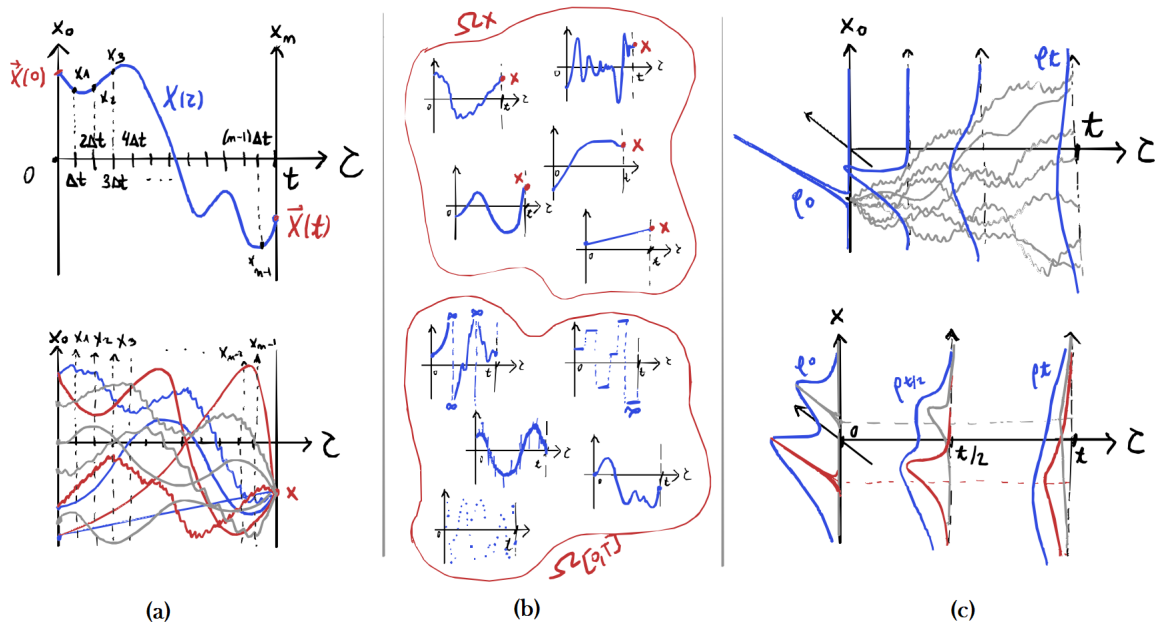


Figure 4: In (a), above, a sketch of the discretization we considered for a trajectory; below, an ensemble of paths ending in a fixed x . Note how sweeping the possible combinations of x_0, x_1, \dots, x_{n-1} essentially exhausts all such paths when $n \rightarrow +\infty$. In (b), above, a sketch of the set Ω_x mentioned in the text, consisting of trajectories ending in x ; below, a sketch of the set of paths $\Omega_{[0,t]} := \prod_{j \in [0,t]} \mathbb{R}^m$ considered in the construction of the Wiener measure, which includes “totally discontinuous” and diverging paths. Despite all, we also found that according to the Wiener measure, continuous paths have measure 1, so the lower sketch does not represent its typicality notion (rather the one above does for $d\mu_{x,t}$). In (c) above, a sketch of the time evolution of the density of a Brownian particle starting at a known location; below, how we can employ the time evolution of a “point source” to get the evolution of a general initial density profile.

(γ) Construction of the Wiener Measure — an Alternative to get Nice Measures in “Infinite Dimensions”

In this section we build a very general limit measure for the path space (now for an arbitrary dimension $m \in \mathbb{N}$ instead of 3) known as the *Wiener measure*. In particular, it allows to find a rigorous version of (121). Now, instead of using explicitly our “trick” above (which will be exemplified in D either-way), here we present an alternative method for limit measures employed by [Reed and Simon \(1975\)](#) in the present context (still, we will find that in its core, our proposed “trick” will tacitly be employed). The essence of the alternative approach relies in using the Riesz-Markov theorem to find a limit measure instead of the Kolmogorov extension theorem. To use the Riesz-Markov theorem, we need that the “field configuration-space”, which in this case is the space of paths (along \mathbb{R}^m and indexed by times in $[0, t]$), has a compact and Hausdorff topology. To achieve it, we will consider the one point compactification of \mathbb{R}^m , adding a point representing the infinity. However, in doing so we will need to accept an even more dramatic tragedy than our loss of norm in §D.2: the compactified \mathbb{R}^m is

not a vector space any-more! Still, just as in *D.4*, the L^2 space of functionals over the space of paths will indeed be a vector space (even a complete Hilbert space!).

That said, lets build the Wiener measures:

- (1). Consider the one-point compactification of \mathbb{R}^m , which we denote by $\overline{\mathbb{R}^m} := \mathbb{R}^m \cup \{\infty\}$ and let $\Omega_{[0,t]} := \prod_{j \in [0,t]} \overline{\mathbb{R}^m}$ be the set of all paths (non-necessarily continuous) γ indexed by a time $\tau \in [0, t]$ and taking values $\gamma(\tau) \in \overline{\mathbb{R}^m}$ (see Figure 4.(b)). We give the path space $\Omega_{[0,t]}$ the product topology, which by Tychonoff's theorem makes $\Omega_{[0,t]}$ a compact and Hausdorff space.
- (2). Let $F : \prod_{j=1}^n \overline{\mathbb{R}^m} \rightarrow \mathbb{R}$ denote a continuous function of n variables. Using it, we can define a function on the path space which only depends on n of the points of a path, $G^{F,t_0,\dots,t_{n-1}} : \Omega_{[0,t]} \rightarrow \mathbb{R}$, via

$$G^{F,t_0,\dots,t_{n-1}}(\gamma) := F(\gamma(t_0), \dots, \gamma(t_{n-1})), \quad (123)$$

for some parameters $t_0 \leq \dots \leq t_{n-1} \in [0, t]$. It is a continuous function (by the continuity of the evaluation map). Then, we can define the set $\mathcal{C}^{fin}(\Omega_{[0,t]})$ as the set of such continuous functionals that depend only on a finite number of entries of the paths and are given by some F . Note that $\mathcal{C}^{fin}(\Omega_{[0,t]})$ is a vector space (even if path-space $\Omega_{[0,t]}$ is not). We equip it with the infinity norm $\|G\|_\infty := \sup_{\gamma \in \Omega_{[0,t]}} |G(\gamma)|$ (which is always finite because G are continuous maps from a compact to \mathbb{R}).

- (3). Fix any $x \in \mathbb{R}^m$, denote $t_n := t$, $x_n := x$ and define the map $L_{x,t} : \mathcal{C}^{fin}(\Omega_{[0,t]}) \rightarrow \mathbb{R}$,

$$L_{x,t}(G^{F,t_0,\dots,t_{n-1}}) := \int_{\mathbb{R}^3} \dots \int_{\mathbb{R}^3} F(x_0, \dots, x_{n-1}) \prod_{j=0}^{n-1} \left(\frac{1}{4\pi(t_{j+1} - t_j)} \right)^{m/2} e^{-\frac{1}{4} \frac{|x_{j+1} - x_j|^2}{|t_{j+1} - t_j|}} dx_j. \quad (124)$$

This is a well-defined⁵³ linear functional acting on $\mathcal{C}^{fin}(\Omega_{[0,t]})$. (Note the formula we obtained in (122)

would be the case for $F(x_0, \dots, x_{n-1}) := e^{\sum_{j=0}^{n-1} \frac{t}{n} V(x_j)} \rho_0(x_0)$ and time nodes $\{0, \frac{t}{n}, 2\frac{t}{n}, \dots, (n-1)\frac{t}{n}\}$.)

- (4). Trivially, if $G^{F,t_0,\dots,t_{n-1}}(\gamma) \geq 0 \forall \gamma \in \Omega_{[0,t]}$ then $L_{x,t}(G^{F,t_0,\dots,t_{n-1}}) \geq 0$, so it is a positive linear functional on $\mathcal{C}^{fin}(\Omega_{[0,t]})$. One can prove that⁵⁴ this implies the operator norm is $\|L_{x,t}\|_{\mathcal{L}(\mathcal{C}^{fin}, \mathbb{R})} = L_{x,t}(1)$, where 1 is the constant function equal to 1 for all paths. But a representative for such a constant function is $F = 1$, and by normalization of the Gaussians, we get $L_{x,t}(1) = 1$. Thus, $L_{x,t}$ is a positive linear functional of norm 1 on $\mathcal{C}^{fin}(\Omega_{[0,t]})$, i.e.

$$|L_{x,t}(G)| \leq \|G(\gamma)\|_\infty \quad \forall G \in \mathcal{C}^{fin}(\Omega_{[0,t]}). \quad (125)$$

- (5). By the Stone-Weierstrass theorem⁵⁵ $\mathcal{C}^{fin}(\Omega_{[0,t]})$ is dense in the set of all continuous functions from $\Omega_{[0,t]}$ to \mathbb{R} , denoted by $\mathcal{C}(\Omega_{[0,t]}, \mathbb{R})$. But then, by the operator extension theorem 3.28 in (Teufel, 2021) (a Hahn-Banach extension theorem), $L_{x,t}$ has a unique positive linear functional extension to all $\mathcal{C}(\Omega_{[0,t]}, \mathbb{R})$ with norm 1 (which we still denote by $L_{x,t}$).

⁵³Given an F of n variables, we could also think of it as a function F' of $N > n$ variables by declaring that F' does not depend on the extra variables (yeah, silly, but we need to consider this to check well-definition of (124)). The extra variables of F' could happen in between those of F , so consider taking (t_0, \dots, t_{n-1}) and (t'_0, \dots, t'_{N-1}) such that the t'_j are equal to the t_k whenever the j -th and k -th are corresponding variables in F and F' . Then, $F'(\gamma(t'_0), \dots, \gamma(t'_{N-1})) = F(\gamma(t_0), \dots, \gamma(t_{n-1})) \forall \gamma \in \Omega_{[0,t]}$ and hence, $G^{F,t_0,\dots,t_{n-1}} = G^{F',t'_0,\dots,t'_{N-1}}$ are the same function (obviously!). The point of this silliness is that because we defined the functional (124) in terms of the representative F , we should check that $L_{x,t}(G^{F,t_0,\dots,t_{n-1}}) = L_{x,t}(G^{F',t'_0,\dots,t'_{N-1}})$. But this is indeed the case because if F' does not depend on some x'_k , we can integrate it out using that for the m dimensional version of (118),

$$\int_{\mathbb{R}^n} K(x_{k-1}, x_k; t_k - t_{k-1}) K(x_k, x_{k+1}; t_{k+1} - t_k) dx_k = K(x_{k-1}, x_{k+1}; t_{k+1} - t_{k-1}).$$

⁵⁴Since for any $G \in \mathcal{C}^{fin}(\Omega_{[0,t]})$ we have $- \|G\|_\infty \leq G \leq \|G\|_\infty$, by positivity and linearity, $-L_{x,t}(1) \|G\|_\infty \leq L_{x,t}(G) \leq L_{x,t}(1) \|G\|_\infty$, or equivalently, $|L_{x,t}(G)| \leq L_{x,t}(1) \|G\|_\infty$, which by definition means $\|L_{x,t}\|_{\mathcal{L}(\mathcal{C}^{fin}, \mathbb{R})} = L_{x,t}(1)$.

⁵⁵Which essentially says that if X is a compact Hausdorff space (in our case $X = \Omega_{[0,t]}$, explaining the need to compactify \mathbb{R}^m), and A is a subalgebra of $\mathcal{C}(X, \mathbb{R})$ (e.g., $\mathcal{C}^{fin}(\Omega_{[0,t]})$) containing a non-zero constant function (e.g., the function 1), then, A is dense in $\mathcal{C}(X, \mathbb{R})$ if and only if it separates points (i.e., for every two points $a, b \in X$ there exists a function $f \in A$ with $f(a) \neq f(b)$) —which in our case holds straightforwardly).

(6). Finally, by the Riesz-Markov theorem⁵⁶ there exists a unique (regular) Borel measure $d\mu_{x,t}$ on $\Omega_{[0,t]}$ with $d\mu_{x,t}(\Omega_{[0,t]}) = 1$ (probability measure) so that

$$L_{x,t}(G) = \int_{\gamma \in \Omega_{[0,t]}} G(\gamma) d\mu_{x,t} \quad \forall G \in \mathcal{C}(\Omega_{[0,t]}, \mathbb{R}). \quad (126)$$

Every such measure $d\mu_{x,t}$ is called a **Wiener measure** on $\Omega_{[0,t]}$.

But now, which ensembles of paths in $\Omega_{[0,t]}$ are Borel and what is their measure according to the Wiener measures? These are important questions because in $\Omega_{[0,t]}$ we have wildly crazy paths, starting from nowhere continuous paths till paths that go to infinity for a third of their trip and then come back. Are they significant? Might it be that all continuous paths have measure zero, so that we are working on an unreasonably big path space?

To answer some of these questions, one can first check that for all $\alpha \in (0, 1]$, the set $\Omega_\alpha \subset \Omega_{[0,t]}$ of α -Hölder continuous paths⁵⁷ is a Borel set (an element in the σ -algebra generated by the open sets). Following this,

Theorem S2. Fix any $x \in \mathbb{R}^m$. Then,

- for all $\alpha \in (0, 1/2)$, $d\mu_{x,t}(\Omega_\alpha) = 1$.
- for all $\alpha \in [1/2, 1]$, $d\mu_{x,t}(\Omega_\alpha) = 0$.

Hence, according to a Wiener measure, almost every path is continuous. That is a relief, because we can then consider only continuous paths in our considerations and still be rigorous. However, almost no path is differentiable! As a consequence, a typical trajectory (i.e., one sampled randomly) following the measure given by any Wiener measure, would be a non-differentiable trajectory. As an example consequence: trajectories for whom the velocity is well-defined and the kinetic energy term in the action is well-defined are almost non-existent under this notion of typicality! So, expressions like (121) are ill-defined. Fortunately, we lumped the kinetic energy part inside the measure and dodged this problem!

(δ) The Feynman-Kac formula

In particular, fixing back $m = 3$, result (126) tells us that for any continuous function $\rho_0 : \mathbb{R}^3 \rightarrow \mathbb{R}$ (as a particular case of an F with a single variable and $t_0 = 0$), now denoting back paths by \vec{X} instead of γ ,

$$\begin{aligned} \int_{\vec{X} \in \Omega_{[0,t]}} \rho_0(\vec{X}(0)) d\mu_{x,t} &= \int_{\mathbb{R}^n} \rho_0(x_0) \left(\frac{1}{4\pi t}\right)^{3/2} e^{-\frac{1}{4} \frac{|x-x_0|^2}{t}} dx_0 = \\ &= \int_{\mathbb{R}^n} \rho_0(x_0) K(x_0, x, t) dx_0. \end{aligned} \quad (127)$$

One can show this formula also holds for any ρ_0 that is measurable on \mathbb{R}^3 and has integrable $\rho_0(\cdot)K(\cdot, x, t)$ (using that such a ρ_0 can be approximated by continuous functions and then applying the dominated convergence theorem). But then (127) holds for all $\rho_0 \in L^2(\mathbb{R}^3)$ and it turns out to be exactly the integral operator of (117). Therefore, we have that rigorously, the path integral formula giving the free time evolution is

$$e^{-tH_0} \rho_0(x) = \int_{\vec{X} \in \Omega_{[0,t]}} \rho_0(\vec{X}(0)) d\mu_{x,t} \quad \forall \rho_0 \in L^2(\mathbb{R}^3). \quad (128)$$

So, we have found that the kernel of the operator e^{-tH_0} , which was $K(x_0, x, t)$ for the Lebesgue measure over \mathbb{R} , according to the Wiener measure is 1 over $\Omega_{[0,t]}$.⁵⁸ The measure took over the role of free propagator. In a similar fashion we can find the rigorification of the full path integral (121), as the limit expression of the also rigorous equation (120):⁵⁹

⁵⁶Essentially it states that given X is a compact Hausdorff space (like our $\Omega_{[0,t]}$, explaining the need to compactify \mathbb{R}^m), for any positive linear functional $L : \mathcal{C}(X, \mathbb{R}) \rightarrow \mathbb{R}$, there exists a unique (positive real regular) Borel measure $d\mu$ on X such that $L(G) = \int_{\gamma \in X} G(\gamma) d\mu \quad \forall G \in \mathcal{C}(X, \mathbb{R})$.

⁵⁷For instance, if $\alpha = 1$, these are the Lipschitz paths, and for any $\alpha \in (0, 1]$ they are uniformly continuous paths.

⁵⁸Analogous to how the Fock “vacuum”, represented by a Gaussian in wavefunctional space was mapped to the function 1 in §D.4.

⁵⁹Note that in these rigorifications we take the whole space of paths over $[0, t]$, instead of only the paths that end in x . The “knowledge” of which is the end-point of the paths that we care about is inside $d\mu_{x,t}$ (note the x dependence). This was evident from the “discrete formulas” like (124), which we employed to build the measure, or (122): the dependence on x was only explicit in the Gaussian that

Theorem S3. (The Feynman-Kac formula) Let V be a real valued potential in $L^2(\mathbb{R}^3) + L^\infty(\mathbb{R}^3)$, $H_0 = -\Delta$ the “free Hamiltonian” and $H = H_0 - V$ the “Hamiltonian” of the heat equation $-\frac{d\rho_t}{dt} = H\rho_t$. Then, for all $\rho_0 \in L^2(\mathbb{R}^3)$,

$$(e^{-tH}\rho_0)(x) = \int_{\vec{X} \in \Omega_{[0,t]}} e^{\int_0^t V(\vec{X}(\tau))d\tau} \rho_0(\vec{X}(0)) d\mu_{x,t}. \quad (129)$$

(This is exactly the ideal formula (121) after we applied the §D.2 trick on (122).)

Proof. We do the proof for compactly supported continuous potentials V . The extension of the result to the rest of potentials, is not as enlightening and is given in Reed and Simon (1975) X.68.

For compactly supported continuous potentials V , we are in the conditions where (120) is rigorous, and hence also (122). As we observed during the construction of the Wiener measure, each element of the sequence in (122), is a particular case of (124) with $F(x_0, \dots, x_{n-1}) = e^{\sum_{j=0}^{n-1} \frac{t}{n} V(x_j)} \rho_0(x_0)$. Hence, by construction of (126), putting back $x_j = \vec{X}(j\frac{t}{n})$ in the potential,

$$e^{-tH}\rho_0(x) = \lim_{n \rightarrow +\infty} \int_{\vec{X} \in \Omega_{[0,t]}} e^{\sum_{j=0}^{n-1} \frac{t}{n} V(\vec{X}(j\frac{t}{n}))} \rho_0(\vec{X}(0)) d\mu_{x,t}, \quad (130)$$

where as always in this appendix, the limit holds in the L^2 sense. This implies in particular that there is a subsequence of (130) for which the convergence to $e^{-tH}\rho_0$ happens almost everywhere in $x \in \mathbb{R}^3$.

We now want to use the dominated convergence theorem on the mentioned subsequence of the limit in (130), in order to “put the limit under the integral” and obtain the desired result (129). We check the ingredients for dominated convergence:

- If \vec{X} is a continuous path, then $t \mapsto V(\vec{X}(t))$ is continuous and hence, by definition of the Riemann integral (which matches with the Lebesgue integral in this case),

$$\lim_{n \rightarrow +\infty} \frac{t}{n} \sum_{j=0}^{n-1} V\left(\vec{X}\left(j\frac{t}{n}\right)\right) = \int_0^t V(\vec{X}(\tau))d\tau. \quad (131)$$

- For every fixed $x \in \mathbb{R}^3$, we found in Theorem 2 that $d\mu_{x,t}$ -almost all paths \vec{X} of $\Omega_{[0,t]}$ are continuous. Thus, we get that point-wise $d\mu_{x,t}$ -almost everywhere in $\Omega_{[0,t]}$,

$$\lim_{n \rightarrow +\infty} \rho_0(\vec{X}(0)) e^{\frac{t}{n} \sum_{j=0}^{n-1} V(\vec{X}(j\frac{t}{n}))} = \rho_0(\vec{X}(0)) e^{\int_0^t V(\vec{X}(\tau))d\tau}. \quad (132)$$

- Moreover, for almost every $x \in \mathbb{R}^3$,

$$\int_{\vec{X} \in \Omega_{[0,t]}} \left| \rho_0(\vec{X}(0)) e^{\frac{t}{n} \sum_{j=0}^{n-1} V(\vec{X}(j\frac{t}{n}))} \right| d\mu_{x,t} \leq e^{t \max|V|} \int_{\vec{X} \in \Omega_{[0,t]}} \left| \rho_0(\vec{X}(0)) \right| d\mu_{x,t} \stackrel{(128)}{=} e^{t \max|V|} (e^{-tH_0}|\rho_0|)(x) < +\infty. \quad (133)$$

certifying there is a dominating L^2 function.

Hence, applying the dominated convergence theorem on (130) to we get that (129) holds for almost all $x \in \mathbb{R}^3$. o.ε.δ.

ended up being fused with the measure. Possibly (should be checked), the subset of $\Omega_{[0,t]}$ consisting of continuous paths ending in a fixed $x \in \mathbb{R}^m$ has measure 1 for $d\mu_x$, while zero for $d\mu_y$ if $y \neq x$. What holds true for sure is that because $\Omega_x := \prod_{j \in [0,t]} \overline{R^m} \times \{x\}$ is also compact and Hausdorff, with the exact same steps as above, we could have done the construction of a measure $d\mu'_{x,t}$ using now only the space of paths ending in x at t . Then we could undoubtedly replace $\Omega_{[0,t]}$ by Ω_x in Theorem 3.