Elements of Environmental Decoherence*

Erich Joos
Rosenweg 2, D-22869 Schenefeld, Germany

Abstract

In this contribution I give an introduction to the essential concepts and mechanisms of decoherence by the environment. The emphasis will be not so much on technical details but rather on conceptual issues and the impact on the interpretation problem of quantum theory.

1 What is decoherence?

Decoherence is the irreversible formation of quantum correlations of a system with its environment. These correlations lead to entirely new properties and behavior compared to that shown by isolated objects.

Whenever we have a product state of two interacting systems - a very special state - the unitary evolution according to the Schrödinger equation will lead to entanglement,

$$\begin{aligned} |\varphi\rangle|\Phi\rangle \quad &\stackrel{t}{\longrightarrow} \quad \sum_{n,m} c_{nm}|\varphi_n\rangle|\Phi_m\rangle \\ &= \quad \sum_n \sqrt{p_n(t)}|\tilde{\varphi}_n(t)\rangle|\tilde{\Phi}_n(t)\rangle. \end{aligned} \tag{1}$$

The rhs of (1) can no longer be written as a single product in the general case. This can also be described by using the Schmidt representation, shown in the second line, where the presence of more than one component is equivalent to the existence of quantum correlations.

If many degrees of freedom are involved in this process, this entanglement will become practically irreversible, except for very special situations. Decoherence is thus a quite normal and, moreover, ubiquitous, quantum mechanical process. Historically, the important observation was that this de-separation of quantum states happens extremely fast for macroscopic objects [17]. The natural environment cannot simply be ignored or treated as a classical background in this case.

Equation (1) shows that there is an intimate connection to the theory of irreversible processes. However, decoherence must not be identified or confused with

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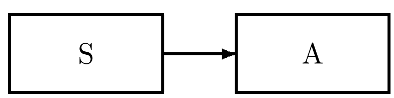
dissipation: decoherence precedes dissipation by acting on a much faster timescale, while requiring initial conditions which are essentially the same as those responsible for the thermodynamic arrow of time [18].

When we consider observations at one of the two systems, we see various consequences of this entanglement. First of all, our considered subsystem will no longer obey a Schrödinger equation, the local dynamics is in general very complicated, but can often be approximated by some sort of master equation (The Schmidt decomposition is directly related to the subsystem density matrices). The most important effect is the disappearance of phase relations (i.e., interference) between certain subspaces of the Hilbert space of the system. Hence the resulting superselection rules can be understood as emerging from a dynamical, approximate and time-directed process. If the coupling to the environment is very strong, the internal dynamics of the system may become slowed down or even frozen. This is now usually called the quantum Zeno effect, which apparently does not occur in our macroscopic world.

The details of the dynamics depend on the kind of coupling between the system we consider and its environment. In many cases – especially in the macroscopic domain – this coupling leads to an evolution similar to a measurement process. Therefore it is appropriate to recall the essential elements of the quantum theory of measurement.

1.1 Dynamical Description of Measurement

The standard description of measurement was laid down by von Neumann already in 1932 [15]. Consider a set of system states $|n\rangle$ which our apparatus is built to discriminate.



Original form of the von Neumann measurement model. Information about the state of the measured system S is transferred to the measuring apparatus A.

For each state $|n\rangle$ we have a corresponding pointer state $|\Phi_n\rangle$ (more precisely, for each "quantum number" n there exists a large set of macrostates $|\Phi_n^{(\alpha)}\rangle$, α describing microscopic degrees of freedom). If the measurement is repeatable or ideal the dynamics of the measurement interaction must look like

$$|n\rangle|\Phi_0\rangle \quad \stackrel{t}{\longrightarrow} \quad |n\rangle|\Phi_n(t)\rangle \ . \tag{2}$$

From linearity we can immediately see what happens for a general initial state of the measured system,

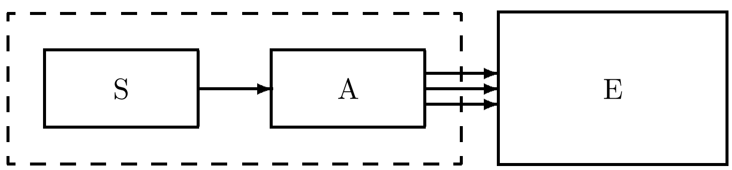
$$\left(\sum_n c_n|n\rangle\right)|\Phi_0\rangle \quad \stackrel{t}{\longrightarrow} \quad \sum_n c_n|n\rangle|\Phi_n(t)\rangle \ . \tag{3}$$

We do not find a certain measurement result, but a superposition. Through unitary evolution, a correlated (and still pure) state results, which contains all possible results as components. Of course such a superposition must not be interpreted as an ensemble. The transition from this superposition to a single component – which is what we observe – constitutes the quantum measurement problem. As long as there is no collapse we have to deal with the whole superposition – and it is well known that a superposition has very different properties compared to any of its components. Quantum correlations are often misinterpreted as (quantum) noise. This is wrong, however: Noise would mean that the considered system is in a certain state, which may be unknown and/or evolve in a complicated way. Such an interpretation is untenable and contradicts all experiments which show the nonlocal features of quantum-correlated (entangled) states.

Von Neumann's treatment, as described so far, is unrealistic since it does not take into account the essential openness of macroscopic objects. This deficiency can easily be remedied by extending the above scheme.

1.2 Classical Properties through Decoherence

If one takes into account that the apparatus A is coupled to its environment E, which also acts like a measurement device, the phase relations are (extremely fast) further dislocalized into the total system – finally the entire universe, according to



Realistic extension of the von Neumann measurement model. Information about the state of the measured system S is transferred to the measuring apparatus A and then very rapidly sent to the environment E. The back-reaction on the (local) system S+A originates entirely from quantum nonlocality.

$$\left(\sum_n c_n|n\rangle|\Phi_n\rangle\right)|E_0\rangle \quad \stackrel{t}{\longrightarrow} \quad \sum_n c_n|n\rangle|\Phi_n\rangle|E_n\rangle. \tag{4}$$

The behavior of system+apparatus is then described by the density matrix

$$\rho_{SA} \approx \sum_n |c_n|^2|n\rangle\langle n| \otimes |\Phi_n\rangle\langle\Phi_n| \qquad \text{if} \qquad \langle E_n|E_m\rangle \approx \delta_{nm} \tag{5}$$

which is identical to that of an ensemble of measurement results $|n\rangle|\Phi_n\rangle$.

Of course, this does not resolve the measurement problem! This density matrix describes only an "improper" ensemble, i.e., with respect to all possible observations at S+A it *appears* that a certain measurement result has been achieved. Again, classical notions like noise or recoil are not appropriate: A acts dynamically on E, but the back-action arises entirely from quantum nonlocality (as long as the measurement is "ideal", that is, (4) is a good approximation). Nevertheless, the system S+A acquires classical behavior, since interference terms are absent with respect to local observations if the above process is irreversible [19, 10].

Needless to say, the interference terms still exist globally in the total (pure) state, although they are unobservable at either system alone – a situation which may be characterized by the statement

The interference terms still exist, but they are not there.[10]

2 Do we need observables?

In most treatments of quantum mechanics the notion of an observable plays a central role. Do observables represent a fundamental concept or can they be derived? If we describe a measurement as a certain kind of interaction, then observables should not be required as an essential ingredient of quantum theory. In a sense this was also done by von Neumann, but not used later very much because of restrictions enforced by the Copenhagen school (e.g., the demand to describe a measurement device in classical terms instead of seeking for a consistent treatment in terms of wave functions).

Two elements are necessary to derive an observable that discriminates certain (orthogonal) system states $|n\rangle$. First, one needs an appropriate interaction which is diagonal in the eigenstates of the measured "observable" and is able to "move the pointer", so that we have as above

$$|n\rangle|\Phi_0\rangle \quad \stackrel{H_{int}}{\longrightarrow} \quad |n\rangle|\Phi_n\rangle \ . \tag{6}$$

This can be achieved by Hamiltonians of the form

$$H_{int} = \sum_n |n\rangle\langle n| \otimes \hat{A}_n \tag{7}$$

with appropriate $\hat{A}_n$ leading to orthogonal pointer states (Note that (6) defines only the eigenbasis of an observable; the eigenvalues represent merely scale factors and are therefore of minor importance). The second condition that must be fulfilled is dynamical stability of pointer states against decoherence, that is, the pointer states must only be passively recognized by the environment according to,

$$|\Phi_n\rangle|E_0\rangle \quad \stackrel{decoherence}{\longrightarrow} \quad |\Phi_n\rangle|E_n\rangle \ . \tag{8}$$

Both conditions must be fulfilled. For example, a measurement device which acts according to (6) would be totally useless, if it were not stable against decoherence: Consider a Schrödinger cat state as pointer state! The *same* basis states $|\Phi_n\rangle$ must be distinguished as dynamically relevant in (6) as well as in (8).*

*This explains *dynamically* why certain observables may "not exist" operationally. For a general discussion of the relation between quantum states and observables see Sect. 2.2 of [5]. Arguments along these lines lead to the conclusion that one should not attribute a fundamental status to the Heisenberg picture – contrary to widespread belief – despite its *phenomenological* equivalence with the Schrödinger picture.

3 Do we need superselection rules?

What is a superselection rule? One way to define a superselection rule is to say, that certain states $|\Psi_1\rangle$, $|\Psi_2\rangle$ are found in nature, but never general superpositions $|\Psi\rangle = \alpha|\Psi_1\rangle + \beta|\Psi_2\rangle$. This means that all observations can be described by a density matrix of the form $\rho = p_1|\Psi_1\rangle\langle\Psi_1| + p_2|\Psi_2\rangle\langle\Psi_2|$. Clearly such a density matrix is exactly what is obtained through decoherence in appropriate situations.

3.1 Approximate superselection rules

There are many examples, where it is hard to find certain superpositions in the real world. The most famous example has been given by Schrödinger: A superposition of a dead and an alive cat

$$|\Psi\rangle = |\text{dead cat}\rangle + |\text{alive cat}\rangle \tag{9}$$

is never observed, contrary to what should be possible according to the superposition principle (and, in fact, *must* necessarily occur according to the Schrödinger equation). Another drastic situation is given by a state like

$$|\Psi\rangle = |\text{cat}\rangle + |\text{dog}\rangle \ . \tag{10}$$

Such a superposition looks truly absurd, but only because we never observe states of this kind! (The obvious objection that one cannot superpose states of "different systems" seems to be inappropriate. For example, nobody hesitates to superpose states with different numbers of particles.) A more down-to-earth example is given by the position of large objects, which are never found in states

$$|\Psi\rangle = |\text{here}\rangle + |\text{there}\rangle \ , \tag{11}$$

with "here" and "there" macroscopically distinct. Under realistic circumstances such objects are always well described by a localized density matrix $\rho(x, x') \approx p(x)\delta(x - x')$. A special case of this localization occurs in molecules (except the very small ones), which show a well-defined spatial structure. The Born-Oppenheimer approximation is not sufficient to explain this fact.

Quite generally we have an approximate superselection rule whenever we describe the dynamics of a dynamical variable by some rate equation (that is, without interference) instead of the Schrödinger equation.

3.2 Exact superselection rules

Strict absence of interference can only be expected for discrete quantities. One important example is electric charge. Can this be understood via decoherence? We know from Maxwell's theory, that every charge carries with itself an associated electric field, so that a superposition of charges may be written in the form [16]

$$\begin{aligned} \sum_q c_q|\Psi_q^{total}\rangle &= \sum_q c_q|\chi_q^{bare}\rangle|\Psi_q^{field}\rangle \\ &= \sum_q c_q|\chi_q^{local}\rangle|\Psi_q^{farfield}\rangle \ . \end{aligned} \tag{12}$$

Since we can only observe the local dressed charge, it has to be described by the density matrix

$$\rho = \sum_q |c_q|^2 |\chi_q^{local}\rangle\langle\chi_q^{local}| \tag{13}$$

If the far fields are orthogonal (distinguishable), coherence would be absent locally. So the question arises: Is the Coulomb field only part of the kinematics (implemented via the Gauss constraint) or does it represent a quantum dynamical degree of freedom so that we have to consider decoherence via a retarded Coulomb field? For an attempt to understand part of the Coulomb field as dynamical see [4].

What do experiments tell us? A superposition of the form (11) can be observed for charged particles (cf. the contribution by Hasselbach[6]). On the other hand, the classical (retarded) Coulomb field would contain information about the path of the charged particle, destroying coherence. The situation does not appear very clear-cut. Hence one essential question remains:

What is the *quantum* physical role of the Coulomb field?

A similar situation arises in quantum gravity, where we can expect that superpositions of different masses (energies) are decohered by the spatial curvature.

Another important "exact" superselection rule forbids superposing states with integer and half-integer spin, for example

$$|\Psi\rangle = |\text{ spin } 1\rangle + |\text{ spin } {}^1/_2\rangle \ , \tag{14}$$

which would transform under a rotation by 2π into

$$|\Psi_{2\pi}\rangle = |\text{ spin } 1\rangle - |\text{ spin } {}^1/_2\rangle \ , \tag{15}$$

clearly a different state because of the different relative phase. If one *demands* that such a rotation should not change anything, such a state must be excluded. This is one standard argument in favor of the "univalence" superselection rule. On the other hand, one *has* observed the sign-change of spin ${}^1/_2$ particles under a (relative) rotation by 2π in *certain* experiments. Hence we are left with two options: Either we view the group SO(3) as the proper rotation group also in quantum theory. Then nothing must change if we rotate the system by an angle of 2π. Hence we can derive this superselection rule from symmetry. But this may merely be a classical prejudice. The other choice is to use SU(2) instead of SO(3) as rotation group. Then we are in need of explaining why those strange superpositions never occur. This last choice amounts to keeping the superposition principle as the fundamental principle of quantum theory. In more technical terms we should then avoid using groups with non-unique ("ray" [¶]) representations, such as SO(3). In supersymmetric theories, bosons and fermions are treated on an equal footing, so it would be natural to superpose their states (what is apparently never done in particle theory).

[¶] The widely used argument that physical states are to be represented by rays, not vectors, in Hilbert space because the phase of a state vector cannot be observed, is misleading. Since relative phases are certainly relevant, one should prefer a vector as a *fundamental* physical state concept, rather than a ray. Rays cannot even be superposed without (implicitly) using vectors.

In a similar manner one could undermine the well-known argument leading from the Galilean symmetry of nonrelativistic quantum mechanics to the mass superselection rule. In this case we could maintain the superposition principle and replace the Galilei group by a larger group. How this can be done is shown by Domenico Giulini[4].

The final open question for this section then is:

Can *all* superselection rules be understood as decoherence effects?

4 Examples

4.1 Localization

The by now standard example of decoherence is the localization of macroscopic objects. Why do macroscopic objects always appear localized in space? Coherence between macroscopically different positions is destroyed *very* rapidly because of the strong influence of scattering processes. The formal description may proceed as follows. Let $|x\rangle$ be the position eigenstate of a macroscopic object, and $|\chi\rangle$ the state of the incoming particle. Following the von Neumann scheme (2), the scattering of such particles off an object located at position x may be written as

$$|x\rangle|\chi\rangle \stackrel{t}{\longrightarrow} |x\rangle|\chi_x\rangle = |x\rangle S_x|\chi\rangle \ , \tag{16}$$

where the scattered state may conveniently be calculated by means of an appropriate S-matrix. For the more general initial state of a wave packet we have then

$$\int d^3x \ \varphi(x)|x\rangle|\chi\rangle \stackrel{t}{\longrightarrow} \int d^3x \ \varphi(x)|x\rangle S_x|\chi\rangle \ . \tag{17}$$

Therefore, the reduced density matrix describing our object changes into

$$\rho(x,x') = \varphi(x)\varphi^*(x') \left\langle \chi|S^\dagger_{x'} S_x|\chi \right\rangle \ . \tag{18}$$

Of course, a single scattering process will usually not resolve a small distance, so in most cases the matrix element on the right-hand side of (18) will be close to unity. If we add the contributions of many scattering processes, an exponential damping of spatial coherence results:

$$\rho(x,x',t) = \rho(x,x',0) \exp\left\{-\Lambda t(x-x')^2\right\} \ . \tag{19}$$

The strength of this effect is described by a single parameter Λ that may be called "localization rate". It is given by

$$\Lambda = \frac{k^2 N v \sigma_{eff}}{V} \ . \tag{20}$$

Here, k is the wave number of the incoming particles, Nv/V the flux, and σ_{eff} is of the order of the total cross section (for details see [10] or Sect. 3.2.1 and Appendix 1 of [5]). Some values of Λ are given in the table.

Localization rate Λ in $\mathrm{cm}^{-2}\mathrm{s}^{-1}$ for three sizes of "dust particles" and various types of scattering processes (from [10]). This quantity measures how fast interference between different positions disappears as a function of distance in the course of time.

	$a = 10^{-3}$cm dust particle	$a = 10^{-5}$cm dust particle	$a = 10^{-6}$cm large molecule
Cosmic background radiation	10^{6}	10^{-6}	10^{-12}
300 K photons	10^{19}	10^{12}	10^{6}
Sunlight (on earth)	10^{21}	10^{17}	10^{13}
Air molecules	10^{36}	10^{32}	10^{30}
Laboratory vacuum (10^3 particles/cm^3)	10^{23}	10^{19}	10^{17}

Most of the numbers in the table are quite large, showing the extremely strong coupling of macroscopic objects, such as dust particles, to their natural environment. Even in intergalactic space, the 3K background radiation cannot simply be neglected.
Hence the main lesson is:

Macroscopic objects are not even approximately isolated.

A consistent unitary description must therefore include the environment and finally the whole universe.*

If we combine this damping of coherence with the "free" Schrödinger dynamics we arrive at an equation of motion for the density matrix that to a good approximation simply adds these two contributions,

$$i\frac{\partial\rho}{\partial t} = [H_{internal}, \rho] + i\frac{\partial\rho}{\partial t}\bigg|_{scatt.} \quad . \tag{21}$$

In the position representation this equation reads in one space dimension

$$i\frac{\partial\rho(x,x',t)}{\partial t} = \frac{1}{2m}\left(\frac{\partial^2}{\partial x'^2} - \frac{\partial^2}{\partial x^2}\right)\rho - i\Lambda(x-x')^2\rho \ . \tag{22}$$

Solutions of this equation can easily be found (see, e.g.[5])

*One of the first stressing the importance of the dynamical coupling of macro-objects to their environment was Dieter Zeh, who wrote in his 1970 Found. Phys. paper [17]: "Since the interactions between macroscopic systems are effective even at astronomical distances, the only 'closed system' is the universe as a whole. ... It is of course very questionable to describe the universe by a wavefunction that obeys a Schrödinger equation. Otherwise, however, there is no inconsistency in measurement, as there is no theory."

This is now more or less commonplace, but this was not the case some 30 years ago, when he sent an earlier version of this paper to the journal Il Nuovo Cimento. I quote from the referee's reply: "The paper is completely senseless. It is clear that the author has not fully understood the problem and the previous contributions in this field." (H.D. Zeh, private communication)

So far this treatment represents *pure* decoherence, following directly the von Neumann scheme. If recoil is added as a next step, we arrive at models including friction, that is, quantum Brownian motion. There are several models for the quantum analogue of Brownian motion, some of which are even older than the first decoherence studies. Early treatments did not, however, draw a distinction between decoherence and friction (decoherence alone does *not* imply friction.). As an example, consider the equation of motion derived by Caldeira and Leggett [2],

$$i\frac{\partial\rho}{\partial t} = [H,\rho] + \frac{\gamma}{2}[x,\{p,\rho\}] - im\gamma k_B T[x,[x,\rho]] \tag{23}$$

which reads for a "free" particle

$$\begin{aligned} i\frac{\partial\rho(x,x',t)}{\partial t} &= \left[\frac{1}{2m}\left(\frac{\partial^2}{\partial x'^2} - \frac{\partial^2}{\partial x^2}\right) - i\Lambda(x-x')^2 \right. \\ &\left. + i\gamma(x-x')\left(\frac{\partial}{\partial x'} - \frac{\partial}{\partial x}\right)\right]\rho(x,x',t)\ , \end{aligned} \tag{24}$$

where γ is the damping constant, and here $\Lambda = m\gamma k_B T$.

If one compares the effectiveness of the two terms representing decoherence and relaxation, one finds that their ratio is given by

$$\frac{\text{decoherence rate}}{\text{relaxation rate}} = mk_BT(\delta x)^2 \propto \left(\frac{\delta x}{\lambda_{th}}\right)^2\ , \tag{25}$$

where λ_{th} denotes the thermal de Broglie wavelength of the considered object. This ratio has for a typical macroscopic situation ($m = 1$g, $T = 300$K, $\delta x = 1$cm) the enormous value of about 10^{40}! This shows that in these cases decoherence is *far more important* than dissipation.

Not only the center-of-mass position of dust particles becomes "classical" via decoherence. The spatial structure of molecules represents another most important example. Consider a simple model of a chiral molecule.

Right- and left-handed versions both have a rather well-defined spatial structure, whereas the ground state is – for symmetry reasons – a superposition of both chiral states. These chiral configurations are usually separated by a tunneling barrier, which is so high that under normal circumstances tunneling is very improbable, as was already shown by Hund in 1929. But this alone does not explain why chiral (and, indeed, most) molecules are never found in energy eigenstates!

In a simplified model with low-lying nearly-degenerate eigenstates $|1\rangle$ and $|2\rangle$, the right- and left-handed configurations may be given by

$$\begin{aligned} |L\rangle &= \frac{1}{\sqrt{2}}(|1\rangle + |2\rangle) \\ |R\rangle &= \frac{1}{\sqrt{2}}(|1\rangle - |2\rangle)\ . \end{aligned} \tag{26}$$

Because the environment recognizes the spatial structure via scattering processes, only chiral states are stable against decoherence,

$$|R,L\rangle|\Phi_0\rangle \stackrel{t}{\longrightarrow} |R,L\rangle|\Phi_{R,L}\rangle \ . \tag{27}$$

The dynamical instability of energy (i.e., parity) eigenstates of molecules represents a typical example of "spontaneous symmetry breaking" induced by decoherence. Additionally, transitions between spatially oriented states are suppressed by the quantum Zeno effect, described below.

4.2 Quantum Zeno Effect

The most dramatic consequence of a strong measurement-like interaction of a system with its environment is the quantum Zeno effect. It has been discovered several times and is also sometimes called "watchdog effect" or "watched pot behavior", although most people now use the term Zeno effect. It is surprising only if one sticks to a classical picture where observing a system and just verifying its state should have no influence on it. Such a prejudice is certainly formed by our everyday experience, where observing things in our surroundings does not change their properties. As is known since the early times of quantum theory, observation can drastically change the observed system.

The essence of the Zeno effect can easily be shown as follows. Consider the "decay" of a system which is initially prepared in the "undecayed" state $|u\rangle$. The probability to find the system undecayed, i.e., in the same state $|u\rangle$ at time t is for small time intervals given by

$$\begin{aligned} P(t) &= |\langle u|\exp(-iHt)|u\rangle|^2 \\ &= 1-(\Delta H)^2t^2+\mathcal{O}(t^4) \end{aligned} \tag{28}$$

with

$$(\Delta H)^2 = \langle u|H^2|u\rangle - \langle u|H|u\rangle^2 \ . \tag{29}$$

If we consider the case of N measurements in the interval $[0,t]$, the non-decay probability is given by

$$P_N(t) \approx \left[1-(\Delta H)^2\left(\frac{t}{N}\right)^2\right]^N > 1-(\Delta H)^2t^2 = P(t) \ . \tag{30}$$

This is always larger than the single-measurement probability given by (28). In the limit of arbitrary dense measurements, the system no longer decays,

$$P_N(t) = 1-(\Delta H)^2\frac{t^2}{N}+\ldots \stackrel{N\to\infty}{\longrightarrow} 1 \ . \tag{31}$$

Hence we find that repeated measurements can completely hinder the natural evolution of a quantum system. Such a result is clearly quite distinct from what is observed for classical systems. Indeed, the paradigmatic example for a classical stochastic process, exponential decay,

$$P(t) = \exp(-\Gamma t) \ , \tag{32}$$

is not influenced by repeated observations, since for N measurements we simply have

$$P_N(t) = \left(\exp\left(-\Gamma\frac{t}{N}\right)\right)^N = \exp(-\Gamma t)\ . \tag{33}$$

So far we have treated the measurement process in our discussion of the Zeno effect in the usual way by assuming a collapse of the system state onto the subspace corresponding to the measurement result. Such a treatment can be extended by employing a von Neumann model for the measurement process, e.g., by coupling a pointer to a two-state system. A simple toy model is given by the Hamiltonian

$$\begin{aligned} H &= H_0 + H_{int} \\ &= V(|1\rangle\langle 2| + |2\rangle\langle 1|) + E|2\rangle\langle 2| + \gamma\hat{p}(|1\rangle\langle 1| - |2\rangle\langle 2|)\ , \end{aligned} \tag{34}$$

where transitions between states $|1\rangle$ and $|2\rangle$ (induced by the "perturbation" V) are monitored by a pointer (coupling constant γ). This model already shows all the typical features mentioned above.

The transition probability starts for small times always quadratically, according to the general result (28). For times, where the pointer resolves the two states, a behavior similar to that found for Markow processes appears: The quadratic time-dependence changes to a linear one. For strong coupling the transitions are suppressed. This clearly shows the dynamical origin of the Zeno effect.

An extension of the above model allows an analysis of the transition from the Zeno effect to master behavior (described by transition *rates* as was first studied in quantum mechanics by Pauli in 1928). It can be shown that for many (micro-)states which are not sufficiently resolved by the environment, Fermi's Golden Rule can be recovered, with transition rates which are no longer reduced by the Zeno effect. Nevertheless, interference between macrostates is suppressed very rapidly [7].

4.3 Decoherence of Fields

In QED we find two (related) situations,

- "Measurement" of charges by fields;
- "Measurement" of fields by charges.

In both cases, the entanglement between charge and field states leads to decoherence as already described above in the discussion of superselection rules, see also [5] and references therein.

In recent quantum optics experiments it is possible to prepare and study superpositions of different classical field states, quantum-mechanically represented by coherent states, for example Schrödinger cat states of the form

$$|\Psi\rangle = N(|\alpha\rangle + |-\alpha\rangle) \tag{35}$$

which can be realized as field states in a cavity. In these experiments (see [1]) decoherence can be turned on gradually by coupling the cavity to a reservoir. Typical decoherence times are in the range of about 100 μs.

For *true* cats the decoherence time is much shorter (in particular, it is *very much* shorter than the lifetime of a cat!). This leads to the appearance of *quantum jumps*, although all underlying processes are smooth in principle since they are governed by the Schrödinger equation.

In experimental situations of this kind we find a gradual transition from a superposition of different decay times (seen in "collapse and revival" experiments) to a local mixture of decay times (leading to "quantum jumps") according to the following scheme.

theory	experiment
superposition of different decay times	collapse and revivals
⇓	⇓
local mixture of different decay times	quantum jumps

4.4 Spacetime and Quantum Gravity

In quantum theories of the gravitational field, no classical spacetime exists at the most fundamental level. Since it is generally assumed that the gravitational field has to be quantized, the question again arises how the corresponding classical properties can be understood.

Genuine quantum effects of gravity are expected to occur for scales of the order of the Planck length $\sqrt{G\hbar/c^3}$. It is therefore often argued that the spacetime structure at larger scales is automatically classical. However, this Planck scale argument is as insufficient as the large mass argument in the evolution of free wave packets. As long as the superposition principle is valid (and even superstring theory leaves this untouched), superpositions of different metrics should occur at any scale.

The central problem can already be demonstrated in a simple Newtonian model[8]. Consider a cube of length L containing a homogeneous gravitational field with a quantum state ψ such that at some initial time $t = 0$

$$|\psi\rangle = c_1|g\rangle + c_2|g'\rangle \ , \tag{36}$$

where g and g' correspond to two different field strengths. A particle with mass m in a state $|\chi\rangle$, which moves through this volume, "measures" the value of g, since its trajectory depends on the acceleration g:

$$|\psi\rangle|\chi^{(0)}\rangle \to c_1|g\rangle|\chi_g(t)\rangle + c_2|g'\rangle|\chi_{g'}(t)\rangle \ . \tag{37}$$

This correlation destroys the coherence between g and g', and the reduced density matrix can be estimated to assume the following form after many such interactions are taken into account:

$$\rho(g, g', t) = \rho(g, g', 0) \exp\left(-\Gamma t(g - g')^2\right) , \tag{38}$$

where

$$\Gamma = nL^4 \left(\frac{\pi m}{2k_B T}\right)^{3/2}$$

for a gas with particle density n and temperature T. For example, air under ordinary conditions, $L = 1$ cm, and $t = 1$ s yields a remaining coherence width of $\Delta g/g \approx 10^{-6}$[8].

Thus, matter does not only tell space to curve but also to behave classically. This is also true in full quantum gravity.

In a fully quantized theory of gravity, for example in the canonical approach described by the Wheeler-deWitt equation,

$$H|\Psi(\Phi, {}^{(3)}\mathcal{G})\rangle = 0 \ , \tag{39}$$

where Φ describes matter and ${}^{(3)}\mathcal{G}$ is the three-metric, everything is contained in the "wave function of the universe" Ψ. Here we encounter new problems: There is neither an external time parameter, nor is there an external observer. How these problems can be tackled is described in Claus Kiefer's contribution[12].

5 Lessons

What insights can be drawn from decoherence studies? It should be emphasized that decoherence derives from a straightforward application of standard quantum theory to realistic situations. It seems to be a historical accident, that the importance of the interaction with the natural environment was overlooked for such a long time. Certainly the still prevailing (partly philosophical) attitudes enforced by the Copenhagen school played a (negative) role here, for example by outlawing a physical analysis of the measurement process in quantum-mechanical terms.

Because of the strong coupling of macroscopic objects, a quantum description of macroscopic objects *requires* the inclusion of the natural environment. A fully unitary quantum theory is only consistent if applied to the whole universe. This does not preclude local phenomenological descriptions. However, their derivation from a universal quantum theory and the interpretation assigned to such descriptions have to be analyzed very carefully.

We have seen that typical classical properties, such as localization in space, are *created* by the environment in an irreversible process, and are therefore not inherent attributes of macroscopic objects. The features of the interaction define *what* is classical by selecting a certain basis in Hilbert space. Hence superselection sectors emerge from the dynamics. In all "classical" situations, the relevant decoherence time is extremely short, so that the smooth Schrödinger dynamics leads to apparent discontinuities like "events", "particles" or "quantum jumps".

There are certain ironies in this situation. *Local* classical properties find their explanation in the *nonlocal* features of quantum states. Usually quantum objects are considered as fragile and easy to disturb, whereas macroscopic objects are viewed as the rock-solid building blocks of empirical reality. However, the opposite is true: macroscopic objects are extremely sensitive and immediately decohered.

On the practical side, decoherence also has its disadvantages. It makes testing alternative theories difficult (more on that below), and it represents a major obstacle for people trying to construct a quantum computer. Building a really big one may well turn out to be as difficult as detecting other Everett worlds!

5.1 Does decoherence solve the measurement problem?

Clearly not. What decoherence tells us, is that certain objects *appear* classical when they are observed. But what is an observation? At some stage, we still have to apply the usual probability rules of quantum theory. These are hidden in density matrices, for example.

5.2 Which interpretations make sense?

One could also ask: what interpretations are left from the many that have been proposed during the decades since the invention of quantum theory? I think, we do not have much of a choice at present*, *if* we restrict ourselves to use only wavefunctions as kinematical concepts (that is, we ignore hidden-variable theories, for example).

There seem to be only the two possibilities either (1) to alter the Schrödinger equation to get something like a "real collapse" [3, 13], or (2) to keep the theory unchanged and try to establish some variant of the Everett interpretation. Both approaches have their pros and cons, some of them are listed in the following table.

Clearly collapse models face the immediate question of how, when and where a collapse takes place. If a collapse occurs before the information enters the consciousness of an observer, one can maintain some kind of psycho-physical parallelism by assuming that what is experienced subjectively is parallel to the physical state of certain objects, e.g., parts of the brain. The last resort is to view consciousness as *causing* collapse, an interpretation which can more or less be traced back to von Neumann. In any case, the collapse happens with a certain probability (and with respect to a certain basis in Hilbert space) and this element of the theory comprises an *additional* axiom.

How would we want to test such theories? One would look for collapse-like deviations from the unitary Schrödinger dynamics. However, similar *apparent* deviations are also produced by decoherence, in particular in the relevant meso- and macroscopic range. So it is hard to discriminate these *true* changes to the Schrödinger equation from the *apparent* deviations brought about by decoherence[9].

Everett interpretations lead into rather similar problems. Instead of specifying the collapse one has to define precisely how the wavefunction is to be split up into branches. Decoherence can help here by selecting certain directions in Hilbert space as dynamically stable (and others as extremely fragile – branches with macroscopic objects in nonclassical states immediately decohere), but the location of the observer in the holistic quantum world is always a decisive ingredient. It must be assumed that what is subjectively experienced is parallel to certain states (observer states) in a certain *component* of the global wave function. The probabilities (frequencies)

*The following owes much to discussions with Dieter Zeh, who finally convinced me that the Everett interpretation *could* perhaps make sense at all.

collapse models	Everett
traditional psycho-physical parallelism: What is perceived is parallel to the observer's physical *state*	new form of psycho-physical parallelism: Subjective perception is parallel to the observer state in a *component* of the universal wave function
probabilities put in by hand	probabilities must also be postulated (existing "derivations" are circular)
problems with relativity	peaceful coexistence with relativity
experimental check: look for collapse-like deviations from the Schrödinger equation ⇓ hard to test because of decoherence	experimental check: look for macroscopic superpositions ⇓ hard to test because of decoherence

we observe in repeated measurements form also an additional axiom [§]. The peaceful coexistence with relativity seems not to pose much problems, since no collapse ever happens and all interactions are local in (high-dimensional) configuration space. But testing Everett means testing the Schrödinger equation in particular with respect to macroscopic superpositions, and this again is precluded by decoherence.

So it seems that both alternatives still have conceptual problems and both are hard to test because of decoherence. We should not be surprised, however, if it finally turned out that we do not know enough about consciousness and its relation to the physical world to solve the quantum mystery [14].

References

[1] Brune, M., Hagley, E., Dreyer, J., Maître, X., Maali, A., Wunderlich, C., Raimond, J.M., Haroche, S. (1996): Observing the Progressive Decoherence of the "Meter" in a Quantum Measurement. Phys. Rev. Lett. **77**, 4887–4890.

[2] Caldeira, A.O., Leggett, A.J. (1983): Path integral approach to quantum Brownian motion. Physica **121A**, 587–616.

[§] There exist several claims in the literature that probabilities can be derived in the Everett interpretation. I think these proofs are circular. Consider a sequence of N measurements on copies of a two-state system, all prepared in the initial state $a|1\rangle + b|2\rangle$. The resulting correlated state contains 2^N components, where each pointer state shows one of the 2^N possible sequences of measurement results (e.g. as a computer printout). But these pointer states are *always the same*, independently of the values of a and b! Only if each branch is given a *weight* involving $|a|^2$ and $|b|^2$ one may recover the correct frequencies. See also [11]. In addition, deep (and partially very old) questions about the meaning of probabilities seem to reappear in the framework of Everett interpretations.

[3] Ghirardi, G.C., Rimini, A. and Weber, T. (1986): Unified dynamics for microscopic and macroscopic systems. Phys. Rev. **D34**, 470–491.

[4] Giulini, D.: States, Symmetries and Superselection, contribution to this volume.

[5] Giulini, D., Joos, E., Kiefer, C., Kupsch, J., Stamatescu, I.-O., Zeh, H.D. (1996): *Decoherence and the Appearance of a Classical World in Quantum Theory* (Springer, Berlin).

[6] Hasselbach. F., Kiesel, H., and Sonnentag, P.: Exploration of the Fundamentals of Quantum Mechanics by Charged Particle Interferometry, contribution to this volume.

[7] Joos, E. (1984): Continuous measurement: Watchdog effect versus golden rule. Phys. Rev. **D29**, 1626–1633.

[8] Joos, E. (1986): Why do we observe a classical spacetime? Phys. Lett. **A116**, 6–8.

[9] Joos, E. (1987): Comment on 'Unified dynamics for microscopic and macroscopic systems', Phys. Rev. **D36**, 3285–3286.

[10] Joos, E., Zeh, H.D. (1985): The emergence of classical properties through interaction with the environment. Z. Phys. **B59**, 223–243.

[11] Kent, A. (1990): Against Many-World Interpretations. Int. J. Mod. Phys. **A5**, 1745–1762. Also available as eprint gr-qc/9703089.

[12] Kiefer, C.: Decoherence in Situations Involving the Gravitational Field, contribution to this volume.

[13] Pearle, P. (1999): Collapse Models. eprint quant-ph/9901077.

[14] Squires, E. (1990): *Conscious Mind in the Physical World* (IOP Publishing, Bristol, Philadelphia).

[15] von Neumann, J. (1932): *Mathematische Grundlagen der Quantentheorie* (Springer, Berlin).

[16] Zeh, H.D.: The Meaning of Decoherence, contribution to this volume.

[17] Zeh, H.D. (1970): On the interpretation of measurement in quantum theory. Found. Phys. **1**, 69–76.

[18] Zeh, H.D. (1999): *The Physical Basis of the Direction of Time.* 3^{rd} edn. (Springer).

[19] Zurek, W.H. (1981): Pointer basis of quantum apparatus: Into what mixture does the wave packet collapse? Phys. Rev. **D24**, 1516–1525.