

$$H_\phi(p_1 \dots q_n, t) := H(p_1 \dots q_n; p_{n+1}(t) \dots q_N(t)). \quad (4.34)$$

Here, particle numbers $1, \dots, n$ are meant to characterize the considered subsystem ‘ ϕ ’, while all others ($n+1, \dots, N$) represent the ‘environment’. Each element of the ensemble would then satisfy another Hamiltonian or Schrödinger equation – in contrast to the assumptions leading to the Liouville or von Neumann equation. Nonetheless, for each element of an ensemble representing incomplete knowledge, the subsystem evolution would be *determined* in this classical case. Neglecting the statistical correlations dynamically by using \hat{P}_{sep} in a master equation would amount to applying the whole resulting ensemble of sub-Hamiltonians (in the forward direction of time) to each individual element of the ensembles of states of the subsystems. However, only under the unstable assumption $\rho = \hat{P}_{\text{classical}}\rho$ (that is, without any entanglement) would the quantum mechanical situation simply be equivalent to the classical one of (4.34), or as in Sect. 3.1.2.

It should be kept in mind, therefore, that the local concepts of relevance, \hat{P}_{sep} , \hat{P}_{local} and $\hat{P}_{\text{classical}}$, appear ‘natural’ only to our classical prejudice. In the unusual situation of controllable entanglement (as in EPR/Bell type experiments), quantum correlations may become relevant by means of the re-localization of superpositions even for local observers. *Dynamical* locality, as described by means of point interactions in field theory, merely warrants the dynamical consistency of these concepts of relevance, or gives rise to the approximate validity of autonomous master equations for $\hat{P}_{\text{local}}\rho$.

General Literature: d’Espagnat 1976, 1983.

4.3 Decoherence

Novel ideas in science are at first completely neglected, then fiercely attacked, and finally regarded as well known.

Konrad Lorenz

In Sect. 3.1 we saw how molecular collisions produce statistical correlations, which describe ‘irrelevant’ information. Although other relevance concepts may also be appropriate for describing irreversible phenomena, the formation of statistical correlations seems to be the most important one in classical description. In a gas, these correlations arise by means of a momentum transfer between molecules, eventually leading to a Maxwell distribution – the distribution of highest entropy for given mean energy if correlations are neglected.

If one specific ‘molecule’ happens to possess macroscopic mass (such as a bullet flying through the gas), its recoil may approximately be neglected in collisions with molecules – except for the resulting friction, whose importance depends on the density of the gas. The bullet may then remain in a non-equilibrium state of almost free motion for some time. On the other hand,

collisions drastically affect the microscopic molecules. Although their states after scattering must strongly depend on the bullet's position at collision time, this dependence cannot be regarded as representing *information* about it if the molecular motions are already chaotic. In contrast, light scattered off the bullet *does* carry information, as we may easily confirm by using our eyes. The reason is that light interacts weakly or coherently with matter, and remains in a state far from equilibrium if absorption can be neglected (see Chap. 2).

The effect of an individual molecule or photon on a macroscopic object may thus be neglected in classical description, but this conclusion has to be radically revised in quantum mechanics. The quantum interaction can be described as an ideal (though uncontrollable) 'measurement' of the bullet's position and shape by the molecule in the sense of von Neumann. If the bullet were initially in a superposition of different positions, as one would have to expect for an object in a generic quantum state, this would lead to an entangled state as in (4.32). In this case, the initial superposition becomes *dislocalized* (it is at *no* place any more). This is called 'decoherence' if the dislocalization is irreversible in practice.³ (Reversible dislocalization of a superposition – such as in a Stern-Gerlach device – may be regarded as 'virtual decoherence'.) It turns out that real decoherence is not only unavoidable for all macroscopic objects, but even the most abundant and most important irreversible process in Nature (Zeh 1970, 1971, 1973, Leggett 1980, Zurek 1981, 1982a, Joos and Zeh 1985).

In general, decoherence is not pure, but accompanied by a distortion of the system under consideration (recoil). For an environmental heat bath this would be required by the *fluctuation–dissipation theorem*, which leads to 'quantum Brownian motion' – a combination of decoherence, dissipation and fluctuation. However, the quantitative relation between these phenomena depends on actual parameters, such as temperature and mass ratios. Since fluctuation and dissipation may so become arbitrarily small, 'ideal' measurements by the environment are appropriate for studying 'pure decoherence' as a genuine quantum phenomenon. Chaotic molecules then contribute to decoherence just as ordered light. Evidently it is the physical effect on the environment that is essential – not any transfer of information. 'Quantum information' is here no more than a misleading renaming of entanglement.

Decoherence is also important for strongly interacting *microscopic* systems, such as individual molecules in a gas, although it is then far from being pure (recoil is essential). Instead of quasi-classical behavior, one now obtains quasi-stochastic dynamics – as successfully *used* in the *Stoßzahlansatz*. Interacting microsystems constituting solids can often be approximated by coupled harmonic oscillators (Caldeira and Leggett 1985). While solutions are then analytically available, they are also known to possess certain pathological properties. In particular, they are non-ergodic.

³ The term decoherence has often been misused in the literature. See Sect. 3.4.3 of Joos et al. (2003) on 'True, False and Fake Decoherence'.

Applying the terminology used in the previous section, decoherence may be understood as the justification of a *specific* $\hat{P}_{\text{semidiag}}$ for a given subsystem by presuming the relevance of locality, as described by the corresponding \hat{P}_{sub} – see (4.31). If this $\hat{P}_{\text{semidiag}}$ turns out to be dynamically valid under all normal circumstances, its eigenspaces characterize ‘quasi-classical’ properties or superselection rules (Zeh 1970, Zurek 1982a). Classical concepts *emerge* approximately in the form of apparent ensembles of narrow wave packets through unavoidable and practically irreversible interaction with the environment. They do not have to be presumed as an independent fundamental ingredient, required for an interpretation of the formalism (as done in the Copenhagen interpretation). From a pragmatic point of view, which does not distinguish between proper and improper mixtures, this would already be sufficient to solve the measurement problem. In a consistent description of reality in terms of wave functions, however, one must assume either a genuine collapse to be *triggered* by decoherence in some way, or appropriately redefine conscious observers within an Everett interpretation (see Sect. 4.6).

The interaction (4.32) was introduced by von Neumann to describe the controllable measurement of a microscopic system ϕ by an appropriate device (with ‘pointer’ states Φ_n). Its fact-like time asymmetry, leading from factorizing to entangled states, could be reversed with sufficient effort if both subsystems were microscopic (‘recoherence’ or ‘erasure of measurement results’). For genuine quantum measurements, the pointer states Φ_n must be macroscopic. They are then ‘measured’ in turn by their uncontrollable environment, and thus become irreversibly quasi-classical. This explains why measurements which lead to macroscopic pointer positions cannot be undone.

It is this universality and unavoidability of entanglement with the natural environment that seems to have been overlooked for the first 50 years of quantum theory. All attempts to describe macroscopic objects quantum mechanically as being isolated, and therefore by means of a Schrödinger equation, were thus doomed to failure – even when including environment-induced dynamical terms that might describe a distortion. Decoherence is different, and extremely efficient, since it does not require an environment that *disturbs* the system. The distortion *of the environment by the system* affects the density matrix of the system, too, because of quantum nonlocality, but on a much shorter time scale than thermal relaxation or dissipation (Joos and Zeh 1985, Zurek 1986).

Some examples of decoherence will now be discussed in detail.

General Literature: Joos et al. 2003, Zeh 2005c, Schloßhauer 2006.

4.3.1 Trajectories

Imagine a two-slit interference experiment with bullets or small dust particles, described by quantum mechanics. Then not only their passage through the slits, but their whole path would be ‘measured’ by scattered molecules or

photons. No interference fringes could ever be observed for such macroscopic objects – even if the resolution of the registration device were fine enough.

In this respect, macroscopic objects are similar to alpha ‘particles’ in a Wilson chamber, which interact strongly with gas molecules by means of their electric charge. For all these objects, their unavoidably arising entanglement with their environment leads to a reduced density matrix that can be represented by an ever-increasing ensemble of narrow wave packets following slightly stochastic trajectories (see also Mott 1929). This result is not restricted to the quantum description of motion in space: propagating wave packets in the configuration space of macroscopic variables may similarly explain their apparent ‘histories’. For spatial motion the argument also demonstrates that the concept of an S -matrix does not apply to macroscopic objects, since it presumes asymptotically free states.

Several very instructive interference experiments have recently been performed with *mesoscopic molecules* that are in the transition region between isolated quantum and classical behavior. Various mechanisms of decoherence, including the emission of thermal radiation from internal molecular degrees of freedom, can be studied for them in detail (Arndt et al. 1999, Hornberger, Hackermüller and Arndt 2005).

For a continuous variable, such as position, decoherence competes with the dispersion of the wave packet that is reversibly described by the Schrödinger equation. Even the scattering rate of photons, atoms, or molecules off small dust particles in intergalactic space suffices to destroy any coherence that would define spreading wave packets for their centers of mass (see Fig. 4.2). If the wavelengths of the abundant scatterers are larger than the width of the wave packet, an otherwise free ‘particle’ is dynamically described by the master equation (Joos and Zeh 1985)

$$i\frac{\partial\rho(x,x',t)}{\partial t} = \frac{1}{2m}\left(\frac{\partial^2}{\partial x'^2} - \frac{\partial^2}{\partial x^2}\right)\rho - i\lambda(x-x')^2\rho. \quad (4.35)$$

It can be *derived* from a universal Schrödinger equation by assuming the dynamical irrelevance of all correlations with the environment after they have formed, and by neglecting recoil (see also Chap. 3 and Appendix 1 of Joos et al. 2003). The coefficient λ is here determined by the rate of scattering and its efficiency in orthogonalizing states of the environment. In the small wavelength limit, a *single* collision is usually sufficient to destroy any coherence beyond the wavelength. The decoherence rate is then simply given by the scattering rate (that is, the product of the flux of environmental particles and the total cross-section). Even the interpretation of the wave mechanical scattering process as consisting of individual collision *events* can be explained by further decoherence of superpositions of different ‘collision times’ in a process that is actually smooth (see Sect. 4.3.6).

So one does not have to *postulate* a fundamental semigroup in order to describe open quantum systems (Sect. 4.4). If the environment forms a heat bath, (4.35) describes the infinite-mass limit of *quantum Brownian motion*

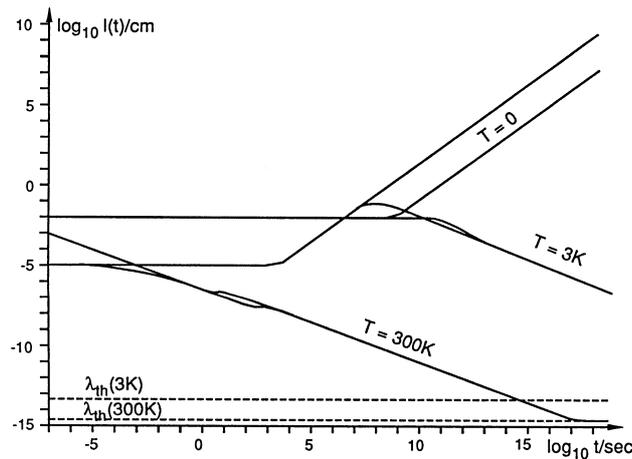


Fig. 4.2. Time dependence of the coherence length $l(t)$ for the center of mass of a small dust grain of 10^{-14} g with radius 10^{-5} cm under continuous measurement by thermal radiation. The six curves represent two initially pure Gaussian wave packets, differing by their initial widths $l(0)$, and three different temperatures T of the radiation. $T = 0$ describes the free dispersion of the wave packet according to the Schrödinger equation, which holds otherwise as an approximation for a limited time only. Scattering of atoms and molecules is in general far more efficient than that of thermal photons – even in intergalactic space. Brownian motion becomes relevant only when the coherence length approaches the de Broglie wavelength λ_{th} . From Joos and Zeh (1985)

(see Caldeira and Leggett 1983, Zurek 1991, Hu, Paz and Zhang 1992, Omnès 1997). This demonstrates that, even for entirely negligible recoil (which would be responsible for noise and friction), there remains an important effect that is based on quantum nonlocality.

Apparent classical properties thus *emerge* from the wave function, and are maintained, by a process that cannot be reversed. In particular, particle aspects (such as tracks in a bubble chamber) arise in the form of macroscopic phenomena (bubbles) which are observed at certain positions in space because of their decoherence. Similarly, the disappearance of interference between *partial waves* in a *Welcher Weg* measurement (Scully, Englert and Walther 1991) does not require any wave-particle ‘complementarity’. Furthermore, no superluminal tunnelling (see Chiao 1998) may occur according to a consistent quantum description, since *all parts* of a wave packet propagate (sub-)luminally, while its group velocity does not represent the propagation of any physical objects.

Master equations for open systems can also be derived by means of the *decoherence functional* (Feynman and Vernon 1963, Mensky 1979). Feynman’s path integral is thereby used as a *tool* for calculating the propagation of a global density matrix, while the environment is again continuously traced out

when getting entangled with the considered system. The intuitive picture of an *ensemble of paths* (representing different *possible trajectories*) is justified only if this superposition of paths decoheres into narrow wave packets. A ‘restriction’ of the path integral by the presence of absorbers (Mensky 2000) would be equivalent to a corresponding reduction of the (total) wave function.

All quasi-classical phenomena, including those representing apparently reversible (friction-free) mechanics, rely conceptually on irreversible decoherence. This requires the continuous production of objective physical entropy (increasing entanglement), which may be macroscopically negligible, but is large in terms of bits. If the quasi-classical trajectories are chaotic, this entropy production may be controlled by the classical Lyapunov exponent (Zurek and Paz 1994, Monteoliva and Paz 2000), even though the entanglement entropy does not require any (initial) uncertainties that would grow *in the direction of calculation*, as assumed in the classical theory of chaos (see Sect. 3.1.2).

General Literature: Joos’s Sect. 3.2 of Joos et al. (2003), Hornberger, Hackermüller, and Arndt (2005).

4.3.2 Molecular Configurations as Robust States

Chirality of molecules, such as right- or left-handed sugar, represents a discrete elementary variable controlled by decoherence. Although a chiral state is described by a certain wave function, it is not an energy eigenstate, which would have to be a parity eigenstate, that is, a symmetric or antisymmetric superposition of both chiralities (see Zeh 1970, Primas 1983, Woolley 1986). The reason is that it is chirality (not parity) that is continuously ‘measured’, for example by scattered air molecules – in analogy to position rather than momentum being measured for a macroscopic ‘mass point’. For sugar molecules under normal conditions, the decoherence time scale is of the order of 10^{-9} s (Joos and Zeh 1985), while the tunneling time between different chirality states is extremely long.

As a consequence, each individual molecule in a bag of sugar retains its chirality, while a parity state – if it had come into existence in a mysterious or expensive way – would almost immediately ‘collapse’ into an apparent ensemble of two chirality states (with equal probabilities). Parity would thus *not be conserved* for sugar molecules, while chirality is always confirmed when measured twice (although it is not a constant of the motion).

This *robustness* against decoherence seems to characterize properties that we usually regard as ‘elements of classical reality’, such as spots on the photographic plate or other ‘pointer states’ of a measurement device. Discrete states may even be protected against otherwise possible transitions (tunneling) by the *quantum Zeno effect*. For continuous variables, the concept of robustness is compatible with a (regular) time dependence according to a master equation, as described in the previous section for the quasi-classical center of mass motion of macroscopic objects. Since entropy production by interaction with

the environment is lowest for a density matrix that is already diagonal in terms of robust states, this property has been called a ‘predictability sieve’, and proposed as a definition of classical states (Zurek, Habib and Paz 1993).

Dynamical robustness is also essential for the physical concept of memory or information storage, such as in DNA, brains or computers. Even ‘states-of-being-conscious’ (see Chap. 1) seem to be quasi-classical in this sense (Tegmark 2000) – at least inasmuch as they are able to communicate. In contrast to such robust properties, which can be assumed to exist regardless of their actual measurement, *potentially* measurable quantities have been called ‘counterfactuals’. Their superpositions, which would themselves describe individual physical states, must not be assumed to describe ensembles of definite (really existing though unknown) properties. Such different concepts of reality (operational or phenomenological versus hypothetical though consistent and economically chosen) can thus be analyzed and understood in terms of decoherence, which is thereby assumed to represent a physical process in a consistent (nonlocal) quantum reality, while elements of phenomenological (classical) reality ‘emerge’ (or become ‘factual’) only under certain environmental conditions. If these conditions may change, such as for microscopic systems under different measurements, the emerging concepts naturally vary between ‘complementary’ modes of description.

Chemists know furthermore that atomic nuclei or strongly bound ions as constituents of *large* molecules have to be described classically (for example as quasi-rigid configurations, which may vibrate or rotate in a time-dependent manner), while the electrons have to be described by stationary or adiabatically comoving wave functions. This asymmetric behavior is often attributed, by means of a Born–Oppenheimer approximation, to their large mass ratio. However, this argument is insufficient, since this approximation applies as well to small molecules that are found in discrete energy eigenstates, which are completely described by stationary wave functions, giving rise to discrete rotational and vibrational energy bands rather than quasi-classical states.

The formation of time-dependent (particle-like) wave packets for the atomic nuclei in large molecules can instead be understood once again by means of decoherence (Joos and Zeh 1985). For example, the positions of nuclei are usually permanently monitored by scattering of lighter molecules that form the environment. But why only the nuclei (or ions), and why not very small molecules? The answer requires a quantitative investigation in each individual case, and the result depends on a delicate balance between internal dynamics and interaction with the environment, whereby the density of states plays a crucial role (Joos 1984). This may then lead approximately to either (a) unitary evolution (including stationary states), (b) a master equation, or (c) freezing of the motion (quantum Zeno effect). Much numerical work remains to be done for such complex systems, while simple ones may be described by an effective master equation, such as (4.35), for example.

General Literature: Joos’s Sect. 3.2.4 of Joos et al. 2003.

4.3.3 Quantum Computers

Digital computers are based on robust binary states, carrying ‘bits’ of information. Even neural networks can be described to some extent by states of cells having ‘fired’ or not, while DNA is based on four different ‘letters’, each one therefore representing two bits. Chiral molecules also represent bits, although they would not be very convenient for information handling.

Just as chiral states may be robust because of their decoherence, so are all macroscopic constituents that are used in classical computers. However, on a microscopic scale there also exist quantum bits (or ‘qubits’), which may occur in all conceivable superpositions of their two basic states. In some cases, such as photon polarizations or spinors that may form spin lattices, they may even be assumed to be isolated from the environment to a good approximation. Such isolated qubits form the essential constituents of quantum computers. Because of their greater variety of possible states (for example spin-up and spin-down in any direction of space), and the possibility of getting entangled, they offer quite novel possibilities for computing (see Shor 1994).

The problem here is that completely isolated systems, required for a unitary evolution, could hardly be manipulated or read as wished for a usable computer. On the other hand, any uncontrollable effect of the collective state of an n -qubit system on the environment would immediately destroy (that is, irreversibly delocalize) the crucial superposition that forms the state of this system as a whole. This vulnerability of quantum computers against decoherence grows exponentially with their size, so that macroscopic quantum computers may have to be excluded by superselection rules, similarly to macroscopic superpositions in general. Superpositions containing a large number of entangled electrons that have been prepared and observed in the laboratory (Mooij et al. 1999, Friedman et al. 2000) are facilitated by ‘freezing out’ most of the degrees of freedom in a degenerate state – in stark contrast to what would be required for the complexity of a quantum computer.

In an attempt to overcome this problem, various *correction codes* have been proposed (see Bouwmeester, Ekert, and Zeilinger 2000). They are conventionally based on some concept of multiple redundancy (an internal kind of back-up), that would have to further enlarge the number of qubits. However, while redundancy may be used as a protection against distortions of the computer by the environment, decoherence is a distortion of the environment by the computer. It can only be corrected for inasmuch as the environment remains controllable – certainly not a very realistic assumption. Usable quantum computers may therefore be excluded in practice for some time to come (see also Haroche and Raimond 1996). It would be quite inconsistent, though, to study the possibility of quantum computers even in principle, while at the same time denying the *reality* of all components of a quantum superposition or wave function – as appropriately emphasized by David Deutsch (1997). Decoherence, too, is the consequence of such an assumption.

In order to give rise to a *classical* computer, each bit would have to be decohered after each calculational step. This would produce precisely the minimum amount of entropy of $k \ln 2$ that was conjectured to be required by Landauer for other reasons (see the end of Sect. 3.3), but then refuted by Bennett in a classical deterministic setting. This entropy production would thus again (have to) be avoided in quantum computers according to the deterministic Schrödinger equation, which is valid only for isolated systems.

General Literature: Bouwmeester, Ekert, and Zeilinger (2000).

4.3.4 Charge Superselection

Gauss' law, $q = (1/4\pi) \int \mathbf{E} \cdot d\mathbf{S}$, tells us that every local electric charge requires a certain flux of electric field lines through a sphere surrounding it at any distance. For a superposition of different charges, one would therefore obtain an entangled quantum state of charges and fields,

$$\begin{aligned} \sum_q c_q \psi_q^{\text{total}} &= \sum_q c_q \chi_q \Psi_q^{\text{field}} = \sum_q c_q \chi_q \Psi_q^{\text{near}} \Psi_q^{\text{far}} \\ &=: \sum_q c_q \chi_q^{\text{dressed}} \Psi_q^{\text{far}}, \end{aligned} \quad (4.36)$$

where χ_q represents the bare charge, while $\Psi_q^{\text{field}} = \Psi_q^{\text{near}} \Psi_q^{\text{far}}$ is the state vector of its correlated electrostatic field, symbolically written as a tensor product of a near field and a far field (see Sect. 2.3). The dressed (physical) charged particle would then be described by a density operator of the form

$$\rho_{\text{local}} = \sum_q |\chi_q^{\text{dressed}}\rangle |c_q|^2 \langle \chi_q^{\text{dressed}}|, \quad (4.37)$$

provided that the states of the far field for different charge q are mutually orthogonal (uniquely distinguishable). The charge is thus decohered by its own Coulomb field, and no charge superselection rule has to be *postulated* (see Giulini, Kiefer and Zeh 1995). The formal decoherence of the bare charge by its near field remains unobservable, since experiments can only be performed with dressed charges.

While this result explains the observed charge superselection rule, one may ask what it means locally. What if an electric charge is accompanied by a negative one at a different place? Or at what distance and on what time scale would the superposition of two different locations of a point charge (such as those of an electron during an interference experiment) be decohered by the quantum state of the corresponding dipole field. A *classical* retarded Coulomb field would contain causal information about the precise path of its source particle. However, interference between different paths of an electron has been demonstrated to exist at least over distances of the order of millimeters (Nicklaus and Hasselbach 1993). This indicates that the Coulomb field contributes

to decoherence only by its monopole component, sufficient to explain charge superselection.

This conclusion can indeed be understood in terms of quantum theory, since photons with diverging wavelength (which may be regarded as representing static fields) cannot distinguish different charge positions – even though the number of such virtual photons would diverge in a Coulomb field. Static dipole (or higher) multipole moments do not possess any far fields. Therefore, only the ‘topological’ Gauss constraint $\partial_\mu F^{\mu 0} = 4\pi j^0$ contributes to the decoherence of the physical particle by the Coulomb field. Any time-dependence (including a retardation) must then be described in terms of transverse photons, represented by the vector potential \mathbf{A} (with $\text{div}\mathbf{A} = 0$ in the Coulomb gauge). In this picture, only the spatial distribution of electric field lines – not their total flux – forms dynamical degrees of freedom that have to be quantized. Charge decoherence has therefore been regarded as ‘kinematical’, although it might as well be assumed to be dynamically *caused* by the retarded field of the (conserved) charge in its past – or equivalently by the advanced field resulting from its future. Note, however, that a kinematical Coulomb constraint is in conflict with the concept of a physical Hilbert space that is spanned by direct products of *local* states.

Dipoles and higher moments (which can define position *differences* for a point charge), can thus be measured by the environment either through emission (or scattering) of transverse (‘real’) photons, or by the irreversible polarization of nearby matter (Kübler and Zeh 1973, Anglin and Zurek 1996). The latter effect has now been experimentally confirmed (Sonnentag and Haselbach 2005). In general, this decoherence is not ‘pure’, but related to energy transfer, although the recoil caused by emission of soft photons may be negligible. The (often virtual) decoherence of individual charged particles *within* solid bodies is discussed in Imry (1997).

The emission of photons would require the charge to be *accelerated*. For example, a *transient dipole* of charge e and maximum distance d , caused by spatially separating opposite charges for a time interval t , requires accelerations a of the order d/t^2 . According to Larmor’s classical formula (see Sect. 2.3), the intensity of radiation is then at least $2e^2a^2/3$. In order to resolve the position difference, the emitted radiation has to consist of photons with energy greater than $\hbar c/d$ (that is, wavelengths smaller than d). The probability that information about the dipole is radiated away by at least one photon is then very small: of order $\alpha Z^2(d/ct)^3$, where α is the fine structure constant and Z the charge number. In more realistic cases, such as interference experiments with electrons, stronger accelerations may occur, but they would in general still cause negligible decoherence.⁴ Decoherence of the position of a charged

⁴ This limitation of the information capacity of an electromagnetic field by its quantum nature must also give rise to an upper bound for the validity of Borel’s argument of Sect. 3.1.2.

particle is therefore dominated by *scattering* of photons, and by interaction with charged or polarizable matter.

The *gravitational field* of a point mass is similar to the Coulomb field of a point charge. Superpositions of different mass should therefore be decohered by the quantum state of the monopole contribution of spatial curvature, and thus give rise to a mass superselection rule. However, superpositions of different energies (hence masses) evidently exist, since they form the time-dependent states of local systems. This situation may not yet be sufficiently understood.

The Coulomb field would vanish globally if the total charge of the Universe were zero (see Giulini, Kiefer and Zeh 1995). This would eliminate the need for a Gauss constraint for the Universe. The gravitational counterpart of this global consequence is the absence of time from a closed Universe in quantized general relativity (the Hamiltonian constraint – see Sect. 6.2).

General Literature: Kiefer’s Sect. 4.1.1 and Giulini’s Chap. 6 of Joos et al. 2003.

4.3.5 Quasi-Classical Fields and Gravity

Not only are the quantum states of charged particles decohered by their fields – quantum states of fields may in turn be decohered by the currents on which they act. In this case, ‘coherent states’, that is, Schrödinger’s time-dependent but dispersion-free Gaussian wave packets for the amplitudes of classical wave modes (eigenmodes of coupled oscillators), have been shown to be robust for similar reasons as electric charges, chiral molecules or the wave packets describing the center of mass motion of quasi-classical objects (Kübler and Zeh 1973, Kiefer 1992, Zurek, Habib and Paz 1993, Habib et al. 1996). This explains why macroscopic states of neutral boson fields appear as *classical fields*, and why superpositions of macroscopically different ‘mean fields’ or different vacua (Sect. 6.1) are never observed.

Coherent harmonic oscillator states, which form states of minimum Heisenberg uncertainty, can be defined (for each wave mode k) as eigenstates $|\alpha_k\rangle$ of the non-Hermitian annihilation (or energy-lowering) operators a_k with their complex eigenvalues α_k (that is, $a_k|\alpha_k\rangle = \alpha_k|\alpha_k\rangle$). These Gaussian wave packets are centered at a time-dependent classical field amplitude $\alpha_k(t) = \alpha_k^0 e^{i\omega t}$, where $\text{Re}(\alpha_k)$ and $\text{Im}(\alpha_k)$ represent the electric and magnetic field strengths, formally equivalent to the position and momentum of a mechanical oscillator. Since the interaction between the field and its charged sources is usually linear in the field operators a_k or a_k^\dagger , these coherent states form an (overcomplete) robust ‘pointer basis’: they create minimal entanglement with their ‘environment’ (that consists here of charged sources that happen to be present).

In contrast to these superpositions of many different photon numbers (or oscillator quantum numbers), *single-photon* states resulting from the decay of different individual atoms (or even the n -photon states resulting from the

decay of a *different number* n of atoms) are unable to interfere with one another, since they are entangled with mutually orthogonal final states of the sources. Two incoherent components of a one-photon state may then appear as ‘different’ photons (using Dirac’s language), although the photons themselves are indistinguishable. A quasi-classical collective state of the source, however, would hardly change (judged in terms of the Hilbert space inner product) when emitting a photon. It is thus able to *produce* the coherent superpositions of different photon numbers discussed above (see also Kiefer 1998).

Although the coherent states behave macroscopically, superpositions of *different* ones, $c_1|\alpha_1\rangle + c_2|\alpha_2\rangle$ (called ‘Schrödinger cat states’), have been produced and maintained for a short time as one-mode laser fields in a cavity (Monroe et al. 1996). These mesoscopic superpositions must decohere, similarly to a Schrödinger cat, although on a time scale that is slow enough to allow this decay of coherence to be monitored as a function of time. In this way, decoherence was for the first time confirmed experimentally as a smooth process in accord with the Schrödinger equation (Davidovich et al. 1996, Brune et al. 1996).

Arguments similar to those used in quantum electrodynamics (QED) apply to *quantum gravity* (Joos 1986, Kiefer 1999 – for applications to quantum cosmology see Chap. 6). Quantum states of matter and geometry must be entangled, and give rise to mutual decoherence. The classical appearance of spacetime geometry is thus no reason *not* to quantize gravity. The beauty of Einstein’s theory can hardly be ranked so much higher than that of Maxwell’s to justify its exemption from quantization. An exactly classical gravitational field interacting with a quantum particle would be incompatible with the uncertainty relations – as has been known since the early Bohr–Einstein debate. The reduced density matrix for the metric must therefore be expected to represent an apparent mixture of different quasi-classical curvature states. Since the observer cannot avoid being correlated to them, spacetime curvature always appears to be classically given – see Sects. 4.6 and 6.2.

Moreover, the entropy and thermal radiation (of all fields) characterizing a black hole or an accelerated Unruh detector (Sects. 5.1 and 5.2) are consequences of the entanglement between relativistic vacua on two half-spaces separated by a horizon (each one forming the environment of the other). This entanglement entropy measures the same type of ‘apparent’ ensemble as the entropy produced according to the master equation (4.35) for a macroscopic mass point. The disappearance of coherence behind a horizon has nonetheless occasionally been regarded as a *fundamental* violation of unitarity, and even as the ultimate source of irreversibility (see Sects. 4.4, 5.1 and 6.2). This appears neither justified nor required (see Kiefer, Müller and Singh 1994, Kiefer 2007, Zeh 2005a).

General Literature: Kiefer’s Chap. 4 of Joos et al. 2003, Kiefer 2004.

4.3.6 Quantum Jumps

Quantum objects are often observed by means of flashes on a scintillation screen or ‘clicks’ of a counter. These macroscopic phenomena are then interpreted as caused by pointlike objects, passing through the observing instrument during a short time interval, while this is in turn understood as evidence for a discontinuous ‘decay event’ (for example, of an atomic nucleus). A *rate equation* for such events is equivalent to a master equation, while a *constant* relative rate would describe exponential decay of the source. Discrete quantum jumps between two energy eigenstates have even been observed for single atoms in a cavity by permanently monitoring their energy, thus enforcing decoherence between energy eigenstates (Nagourney, Sandberg and Dehmelt 1986, Sauter et al. 1986, Pegg, Loudon and Knight 1986, Gleyzes et al. 2006). Therefore, formal creation and annihilation operators are often misunderstood as defining discrete events, even though they occur in a Hamiltonian that constitutes a Schrödinger equation.

This Schrödinger equation would describe a state vector that smoothly develops components with different particle numbers, or a wave function that leaks out of an unstable system (such as a quantum ‘particle’ in a potential well). This contrast between discrete events and the Schrödinger equation is clearly the empirical root of the probability interpretation of the wave function in terms of events and particles. A wave function can exponentially decay only in a limited region of space (for example within an expanding sphere for a limited time – see Sect. 4.5). This wave function is a *superposition* rather than an ensemble of different decay times. Their interference and the dispersion of the corresponding outgoing wave lead to deviations from an exponential decay law. Although these deviations are too small to be observed for decay into infinite space, interference between different decay times has often been confirmed in other situations, not least as ‘coherent state vector revival’ for photons emitted into cavities with reflecting walls (Rempe, Walther and Klein 1987).

In Sect. 4.3.1, the appearance of particles following tracks in a cloud chamber has been explained in terms of an apparent ensemble of narrow wave packets arising by means of decoherence. Similar arguments may as well explain apparently discrete events. Even if quantum objects remain isolated before being detected, they would be decohered in the detector – usually on a very short time scale. Therefore, the same decoherence that describes localization in space also explains localization in time. Jumps between discrete energy levels, observed under continuous measurement, represent apparently discrete ‘decay histories’, which can be explained by Mott-type quantum correlations between successive measurements of short but finite individual duration (including the decoherence of their outcomes). Neither particles nor genuine quantum jumps are required as *fundamental* concepts in quantum theory (Zeh 1993, Paz and Zurek 1999). Whenever decay fragments (or the decaying object) interact appropriately with their environment, interference between two partial waves

describing a decayed state and a not yet decayed state disappears on a very short (though finite) decoherence time scale, thus giving rise to an apparent *ensemble* of decay times. This time scale is in general much shorter than the time resolution of measurements.

If the decay status is thus permanently ‘monitored’ by the environment, a set of identical decaying objects is thus more appropriately described by a rate equation than by a Schrödinger equation (Sect. 4.5). This rate equation leads to an exact exponential law, since it excludes any interference between different decay times. Similarly, decay products emitted in superpositions of sufficiently different energies are absorbed into mutually orthogonal final states of the environment. Microscopic systems with their discrete energy levels must therefore decohere into eigenstates of their own Hamiltonians. This explains why the atomic world is characterized by stationary states, and von Neumann spoke of an *Eingriff* (intervention) required for their change.

So it seems that this situation of continuously monitored decay has led to the myth of quantum theory as a stochastic theory for fundamental *quantum events* (see Jadczyk 1995). Bohr (1928) remarked that “the essence” (of quantum theory) “may be expressed in the so-called quantum postulate, which attributes to *any atomic process* an essential discontinuity, or rather individuality . . .” (my italics). This statement is in conflict with many microscopic and mesoscopic quantum phenomena that have since then been observed. Heisenberg and Pauli similarly emphasized their preference for matrix mechanics because of its (evidently misleading) superiority in describing discontinuities. Ole Ulfbeck and Aage Bohr (2001) recently emphasized the unpredictable occurrence of ‘clicks in the counter’, while denying the existence of any quantum events in the source that would precede them. This comes close to the consequences of decoherence, but rather than taking into account entanglement with the environment the authors conclude that “the wave function then loses its meaning”. According to the decoherence theory, the underlying entanglement processes are always smooth, and described by a Schrödinger equation. The short decoherence time scales lead to the impression of quantum *jumps* between energy eigenstates, for example, while narrow wave packets are interpreted as particles or classical variables (even though the certainty of classical properties has to be restricted by the uncertainty relations in order to comply with the Fourier theorem).

While the description of all physical phenomena in terms of time-dependent entangled wave functions now appears as a consistent picture, an important question remains: how should the probabilities, which were required to justify the concept of a density matrix in Sect. 4.2, be understood if they are *not* probabilities for quantum jumps or for the occurrence of measurement results in the form of fundamental ‘events’. This discussion will be resumed in Sect. 4.6.

General Literature: Joos’s Sect. 3.4.1 of Joos et al. 2003.