**Quantum Information and Locality**

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**1. Introduction. What makes quantum information different: non-orthogonality and entanglement**

The surprising aspects of quantum information are due to two distinctly non-classical features of the quantum world: first, different quantum states need not be orthogonal and, second, quantum states may be entangled. Of these two significant non-classical features non-orthogonality has to some extent become familiar and may therefore be called the less spectacular one; we shall only briefly review its consequences for information storage and transfer, in this introductory section. By contrast, the conceptual implications of entanglement are amazing and still incompletely understood, in spite of several decades of debate. The greater part of this chapter (from section 1 onward) is therefore devoted to an attempt to better understand the significance of entanglement, in particular for the basic physical concepts of “particle” and “localized physical system”. It will turn out, so we shall argue, that the latter concepts have only limited applicability and that quantum mechanics is accordingly best seen as not belonging to the category of space-time theories, in which physical quantities can be defined as functions on space-time points. The resulting picture of the physical world is relevant for our understanding of the precise way in which quantum theory is non-local and sheds light on how quantum information processes can succeed in being more efficient than their classical counterparts.

Both in quantum and classical physics physical information is conveyed by messages that are encoded in states of physical systems; these states can be thought of as letters in an alphabet. Information transfer consists in the preparation of such messages by a sender and their subsequent reproduction (possibly in states of other systems) at the location of a receiver.

This characterization may seem to entail that the prospects of efficient information storage and transfer are generally worse in quantum theory than in classical physics. The reason is that in quantum theory different states of physical systems, even if maximally complete (pure states), may “overlap”: they need not be represented by orthogonal vectors in Hilbert space. This non-orthogonality stands in the way of unequivocally distinguishing between states and so between the letters of the alphabet that is being used. Consider, for example, a message written in a binary code using states  and , with ; attempts to decode the message via measurements will inevitably lead to ambiguities. By contrast to the case of (ideal) measurements on classical states or on orthogonal quantum states, in which it is always possible to completely retrieve the original message, performing yes-no measurements represented by the projection operators  and  introduces a probability  of incorrect classification.

More sophisticated measurement schemes exist that do enable classification without error (Dieks 1988; Jaeger 1995). However, the price to be paid is that in these schemes some measurements will have outcomes that lead to no classification at all. The best one can achieve this way is that a fraction  of all bits remains unclassified, so that non-orthogonality again leads to a loss of information.

The possibility of distinct though non-orthogonal states has the further consequence that statistical mixtures of quantum states cannot be uniquely decomposed into pure components. A statistical (impure) state , prepared by probabilistically mixing an arbitrary number of non-orthogonal states  with associated probabilities  is the same mathematical entity as (and empirically indistinguishable from) the mixture of orthogonal states  with probabilities , , where the  and  are the eigenstates and eigenvalues of the density operator  () and where *n* denotes the number of dimensions of the system’s Hilbert space. From an informational viewpoint, the latter mixture (of at most *n* orthogonal and therefore perfectly distinguishable quantum states) is equivalent to a classical probability distribution with probabilities  over *n* distinct classical states. Therefore, the amount of information that can be extracted from is generally less than the (classical) information needed to specify the mixing procedure, in which the number of mixed states may be much larger than *n*.

This difference is made precise by the *noiseless quantum coding theorem* (Schumacher 1995). Essentially, this theorem tells us that the amount of information contained in a mixed state is given by the *von Neumann entropy* , . This von Neumann entropy of is smaller than (or at most equal to) the classical *Shannon entropy* of the probability distribution *pi* that determined the mixing process from which  originated. The latter entropy quantifies the amount of classical information needed to describe the mixing (where we take the different components  as distinct classical possibilities). So instead of the classical result (Shannon 1948) that the Shannon entropy determines the number of bits needed to specify a state, we find that for quantum messaging the use of the smaller von Neumann entropy is appropriate.

This brief sketch illustrates how the *concept* of information does not change by the transition to quantum theory: as in the classical case, we are dealing with messages encoded in states of physical systems and with their transport through a physical channel. But the differences between classical systems and quantum systems, mirrored in the physical properties of their states, are responsible for differences in the *amounts* of information that can be stored and transferred.

The impression that quantum messaging is at best equally powerful as its classical predecessor would be incorrect, however. There are also situations in which quantum mechanics predicts more effective information transfer. The essential ingredient of these new quantum possibilities is the consideration of *composite* systems in *entangled states*. Entangled states represent systems that are correlated in a stronger way than deemed possible in classical physics, and enable instantaneous non-local influences between systems. This typical quantum non-locality underlies revolutionary new quantum information protocols like teleportation. At least, this is the common view on the meaning of entanglement; we shall analyze it, and also criticize it, below.

**2. Particle individuality and labels, classical and quantum**

The conceptual landscape of classical physics is dominated by the notion of a localized individual system, with a *particle* as its paradigm case. Classical particles can be characterized by individuating sets of properties, like position, momentum, mass and electric charge. Different classical particles may possess the same masses, charges and momenta, but they cannot find themselves at the same position since repulsive forces are assumed to become increasingly strong when mutual distances become smaller. Consequently, two distinct particles differ in the values of at least one of their physical characteristics, namely their localization.

Within the framework of Newtonian physics, in which absolute space legitimates the notion of absolute position as a physical property, this guarantees that a version of Leibniz’s principle of the identity of indistinguishable objects is respected: Newtonian particles are distinguishable from each other on the basis of their complete sets of physical properties and this distinguishability grounds their individuality. Within a relational space-time framework a weakened Leibnizean notion, namely that of weak discernibility, may be invoked to ground individuality (in cases like that of Black’s spheres; see Dieks and Versteegh 2008 and references contained therein for further discussion). The general message is that in classical physics the specification, counting and labeling of particles is done on the basis of physical distinctions, position being the most important.

Numerical particle labels () accordingly correspond to such sets of physical properties; they have no content beyond this and function as conventionally chosen names. Thus, although it is true that a numerical labeling induces an ordering on sets of particles, this ordering is without physical significance and may be replaced by any other, permuted ordering: labels can be permuted in the equations without any physical consequence. It is therefore possible to switch over to a completely symmetrical description, in which all states that follow from each other by permutations of the particle labels are combined and the resulting complex is taken as the representation of *one* physical state (the “Z-star” introduced by Ehrenfest (1909)). For example, in the simple case of two particles of the same kind (same mass, charge, and other kind-specific physical properties) −two electrons, say− we could have that one particle is at position  with momentum  and the other at position , with momentum . This gives rise to two points in “phase space” (with axes , , ,  representing the momenta and positions of the two particles, respectively): , , ,  and , , , . But these two points correspond to the exact same physical situation, so that one is entitled to view the Z-star comprising *both* as representative of this single physical situation. This yields a description that is symmetrical in the particle labels and avoids an empirically empty “numerical identity” of the particles (alternatively, one could switch to a “reduced phase space”, by “quotienting out” the equivalence generated by the 1↔2 permutation).

The essential point is that classical systems are individuated by co-instantiated sets of physical properties and that labels possess the status of conventional names for such sets −they contain no information about the identity and individuality of particles beyond this. That physical properties are co-instantiated in fixed combinations is a basic principle of classical physics, with law-like status: physical properties only occur in specific bundles (like mass, position and momentum, etc.) and not in isolation. These bundles define the physical individuals (Lombardi and Dieks 2015).

Quantum mechanics *prima facie* respects a similar “bundle principle”. Indeed, quantum systems like electrons are characterized by fixed sets of physical quantities: e.g., charge, mass, spin and position. Compared to classical physics there is the complication that not all these quantum quantities can simultaneously be definite-valued, as a consequence of complementarity. (The characterization of quantum particles that we give here is not interpretation independent: for example, in interpretations of the Bohm-type it is a first principle that the world consists of particles with always definite positions, contrary to what we assume here. Our analysis proceeds within a broadly construed no-hidden-variables approach). Thus, position and momentum cannot possess sharp values simultaneously (although it is still possible to have both quantities in the same list of particle properties if one allows a certain unsharpness in their values −or “latitudes”, as Bohr called it (Dieks 2016)). This incompatibility between sharp values of certain quantum quantities is taken into account through the use of non-commutative operator algebras instead of numerical quantities for the physical properties, plus the use of quantum states (vectors in Hilbert space) instead of classical states like . In this way we can generalize the classical idea of a set of physical properties whose values together define an individual system, by specifying an algebra of “observables” (hermitian operators on a Hilbert space), together with a pure quantum state. In such a state the observable represented by the projection on the state, plus all observables commuting with this projection operator, are assigned definite values. This restricted value ascription is the best quantum mechanics can achieve: a complete set of quantum properties that takes the place of the classical property sets. Observables of which the state is not an eigenstate are characterized by probability distributions specifying the probabilities of possible outcomes in a measurement of the observables in question.

Modulo this complication of complementarity and non-definiteness in the numerical sense, the situation may seem not too different from the classical one: for example, instead of the values of position, momentum, and angular momentum characterizing a spinning classical particle we may now write down the corresponding position, momentum and spin operators that characterize the system, e.g. an electron, and we can use a state like  to characterize its space-time characteristics through  plus its spin through .

Summing up, quantum particles can be characterized by fixed algebras of operators plus pure quantum states. (*Mixed* states do not characterize a single actual system when they are proper mixtures. Why improper mixtures, coming from entangled states, do not individuate particles is explained below). The algebra plus the state give us a complete set of quantum properties. The difference with classical mechanics thus appears to be solely in the replacement of definite numerical values of quantities with probability distributions determined by the quantum states. However, we shall see that things are not this easy, and that there are reasons to doubt the general applicability of the concept of a particle.

In the case of more than one particle quantum theory works with labels that are meant to index the particles, in the same spirit as in classical physics. As before, we expect and require that labels only individuate by proxy, via physical properties or states. As we shall see, in the quantum case it is more difficult than before to put this simple principle into practice.

First, there is a technical complication due to the fact that the quantum mechanical state space is a vector space, in which the superposition principle holds. The latter has the consequence that the natural way of forming symmetrical states is different from the Z-star procedure explained above. Instead of the *collection* of permuted states, one now takes their *superposition*. In the case of particles of the same sort this is precisely what the (anti-) symmetrization postulates demand: instead of considering product states of the form  (for a situation in which particle 1 is located to the left, particle 2 to the right) or  (the same physical situation, but with the particle labels permuted), we should write down (anti-) symmetrical states of the form

 (1)

The state (1) is the quantum counterpart of Ehrenfest’s symmetrical Z-star for two particles with the same intrinsic properties. If we look at the labels occurring in Eq.(1) we see that they are symmetrically distributed over the two terms in  −so if we were to associate physical systems with these labels we would be forced to conclude that the two systems of (1) are in exactly the same state and therefore cannot be individuated by physical properties. The same conclusion follows in more technical fashion by taking “partial traces”: tracing out over the parts of the total wave function labeled by 2 we obtain a mixed state for “system 1”, namely ; and exactly the same state for “system 2” by tracing out over 1.

We have already observed in the classical case, however, that it would be wrong to uncritically rely on the referential power of labels: we should associate physical systems not with labels *per se* but with physical properties. Just as the symmetry of the classical Z-star does not imply that no individual particles, each one with its own state, exist, the mere symmetry of  in (1) does not compel us to abandon the notion of individual quantum particles. We should ask ourselves whether it is possible to individuate such particles by properties or states, rather than focusing on labels. In other words: we should be prepared to relax the one-to-one link between the labels that occur in the formalism (mathematically introduced as indices of the factor Hilbert spaces in the total tensor product space) and physical systems. The analogy with the classical case in fact suggests that the state  of Eq.(1) corresponds to a situation with one particle located Left and one particle located Right, in particular if  and  stand for spatially widely separated wave packets that do not overlap.

Before embarking on the question of whether such an interpretation is justified, the following general remark is in order. The fact that a *superposition* of quantum states appears in Eq.(1), instead of a simple juxtaposition of possibilities, has an immediate non-classical consequence: the expectation value of an arbitrary operator in the state , , will contain an interference term . If this interference term is non-zero, i.e., when we are dealing with operators that have non-vanishing matrix elements between  and  and thus “connect” *L* and *R*, quantum mechanics predicts results for measurements of the physical quantity represented by  that are of the same type as the interference effects found in double slit experiments. So even if an interpretation of (1) in terms of “one particle to the left and one particle to the right” is going to prove tenable, this will not mean a return to a *classical* particle picture: the systems that we are trying to identify are and remain quantum (Ladyman, Linnebo and Bigaj (2013) discuss some more details of the interference between the two “branches” in Eq.(1) and its empirical consequences in collision processes; they also point out that such interference terms can easily vanish, namely when there are additional physical features that distinguish between  and ).

States of the type (1) are not product states, and this is usually taken to be the defining characteristic of entanglement. Furthermore, it is standard wisdom that identical particles in entangled states −electrons, say− do not possess pure states of their own, and that their mixed states are all the same. Therefore, physical individuation of component systems in such states seems impossible from the outset. But as we have noted, this standard account relies on the supposition that if there are individual systems the *labels* occurring in the many-particles formalism will have to refer to them; and it is exactly this supposition that we are going to dispute.

To make the idea of the argument against the signiﬁcance of “particle labels” clear, consider the EPR-Bohm state, notorious from the locality debates. The EPR experiment (in its modern Bohm version) is standardly discussed as being about two electrons at a large distance from each other, whose spins are correlated because the combination of their spins is described by the singlet state . The spatial part of the wave function is often not written down explicitly, but it obviously must be considered as well in order to make contact with the locality question. The total state including this spatial part has the form

 (2)

where  and  are states localized on the left and right, respectively, at a large distance from each other. Note that the spatial part of  is symmetric in the “particle labels” 1 and 2, as required by the antisymmetrization postulate (the total state must be antisymmetric, since we are dealing with fermions).

Now, if we use the state of Eq.(2) and also think that the labels 1 and 2 refer to our two particles, we have to conclude that there is neither a left nor a right electron. The spatial states associated with both 1 and 2 are exactly the same, namely , so that the two particles defined this way would be “evenly distributed” between left and right. This means that the way the EPR case is standardly understood, as being about two localized particles at a large distance from each other, is at odds with the interpretation of the indices 1 and 2 as particle labels.

In the literature on EPR the state is sometimes given in a different form, namely as

 (3)

in which the spatial part is a simple product state. This state is incorrect, because it violates the antisymmetrization postulate −there are empirical differences between the states (2) and (3), because of the possibility of interference between the branches, as we have noted before. There can be no question that (2) is correct and (3) is not. But it is nevertheless understandable that (3) is sometimes written down, because this form of the state lends itself easily to the desired intuitive interpretation, namely that there is one single particle (particle 1) at the left, and another (particle 2) at the right. In Eq.(3) there is one-to-one link between the labels and the physical location properties *L* and *R*, and this justifies the use of the labels as referring to the Left and Right regions, respectively. (A picture with particle 1 at *L* and particle 2 at *R* is still problematic, though: although the labels 1 and 2 correlate to positions via  and , respectively, there are no complete sets of quantum properties correlated to them. We shall discuss this issue extensively in the next section.)

However, for the sake of correctness we must work with the state (2) instead of (3). In this correct state there is no correspondence at all between individuating physical particle properties and the labels 1 and 2. This poses a dilemma: either we have to conclude that there are no localized systems in EPR-like situations, or we have to look for another way of referring to particles than by the indices occurring in the many-particles formalism.

Apparently, we have hit upon an important difference between classical and quantum physics. In classical physics we can always individuate particles, or more generally physical objects, by unique sets of values of physical quantities. This justifies the use of particle labels, which serve as an abbreviation of such sets. By contrast, in quantum theory it is at least unclear whether objects can be referred to in this way −this might already raise doubt about the general appropriateness of concepts like “particle” and “localized object” in the quantum context.

**3. Identical quantum particles**

If it were the case that the “particle labels” in an EPR-like situation constituted our only basis for speaking about physical systems, we would be forced to accept (among other things) that these systems are not localized in small regions in space: they would be “half in L and half in R”. A measurement on the left wing of the EPR experiment could in this case perhaps equally affect particle 1 and particle 2. This might be considered to possess a positive side: it could perhaps be taken as the starting point for an explanation of the non-local quantum effects that lie at the basis of violations of Bell inequalities, and also of the effectiveness of quantum information processes. If the physical systems that are involved in such processes are non-local to start with, non-locality in their behavior should hardly surprise us. But the price to pay for attempts to use this strategy in this exact form would be high. We would be compelled to accept that all identical particles, electrons for example, would be smeared out over the universe. And, importantly, we would be unable to recover the classical particle picture in the classical limit of quantum theory. In the classical limit it should surely be possible to think in terms of localized particles that can be labeled on the basis of their positions and trajectories −but this means that these classical labels cannot be in one-to-one correspondence with the indices occurring in the quantum formalism (Dieks and Lubberdink 2011). So there must be a different and better way of identifying quantum particles if we wish to have a smooth transition from quantum to classical. When we have found how to do this, we should look at the explanation of non-local effects afresh.

Fortunately, an alternative way of identifying quantum systems is not difficult to find. The crux is in our earlier observation that particle labels have no physical content in themselves and only obtain such content via association with physical quantities. These physical quantities, and physical states, are consequently primary; the labels are secondary auxiliary tools. As Ghirardi, Marinatto and Weber (2002) point out (see also the discussion by Ladyman, Linnebo and Bigaj (2013) and by Caulton (2014)), in the case of particles of the same kind (identical particles) it is accordingly necessary to allow only physical quantities that are symmetrical in the labels. Indeed, think back of Ehrenfest’s Z-star for two classical particles: there, we should count ,  and ,  as *one* physical situation, which will automatically happen if we symmetrize the physical quantities with which we characterize the situation. Analogously, instead of considering particle properties represented by projection operators of the form , with  the unity operator in Hilbert space 2, we should in the quantum case focus on projection operators like

 (4)

The expectation value of the projector in (4), in a properly (anti)symmetrized state, gives the probability of finding at least one of the identical particles in the state onto which *P* projects. The last term in (4) can be left out without consequences in the case of fermions; it has been added to allow for the possibility that the two particles are found in exactly the same state, which could occur when we are dealing with bosons. Let us stress that the analogy with the classical case is not the only reason for symmetrization of the observables: consistency of the formalism requires it as well. For example, a non-symmetrical Hamiltonian would steer the state out of its subspace of (anti)symmetrical states and would thus create a conflict with the (anti)symmetrization postulate.

As Ghirardi, Marinatto and Weber (2002) observe, it is in this way possible to assign a complete set of quantum properties to components of a multi-particles system, even in symmetrical or antisymmetrical states, if the following condition is satisfied: the total (entangled) state should be obtained by symmetrizing or antisymmetrizing a *factorized* state. Such states are eigenvectors of symmetrical projection operators of the type (4). Making use of these sets of quantum properties, we may identify subsystems via their *states* instead of via labels.

This identification of subsystems is in line with the laboratory practice of physics; it also respects the classical limit (Dieks and Lubberdink 2011; Dieks 2014). To see the essential idea in a simple example, consider the case of two fermions in state

 (5)

with  and  two non-overlapping wavefunctions at a large distance from each other. The property assignment by means of the projection operators  and  leads to the conclusion that we have one particle characterized by  and one particle characterized by , respectively (we have left out the last term of (4) in the projection operators, as this term automatically vanishes in the case of fermions). The state (5) represents therefore one particle at the left and one at the right, in spite of the fact that the labels 1 and 2 are evenly distributed over *L* and *R*. In the classical limit this would correspond exactly to what we expect: particles, e.g. electrons, represented by narrow wave packets, following approximately classical trajectories (narrow wave packets spread out, and will therefore only be able to follow approximately classical trajectories for a limited time; for the classical limit description to be applicable additional conditions must be fulfilled, like the presence of decohering mechanisms, see Rosaler 2016). These classical objects can of course be labeled, on the basis of their unique positions and trajectories, but the thus defined labels will not coincide with the indices 1 and 2 occurring in the quantum state −the latter remain evenly distributed over the wave packets even in the classical limit.

We have glossed over an important point, namely that the decomposition of states like (1) generally is not unique, because of the equality of the coefficients appearing in front of the terms in the superposition (“degeneracy”). So alternative descriptions, in addition to the one in terms of  and , are possible. To make the description unique the symmetry needs to be broken; it is plausible that such a symmetry breaking will be realized when the two-particle system enters into interaction with a macroscopic environment, via position-dependent interactions. This issue is important for non-collapse interpretations of quantum mechanics, like the many-worlds interpretation or the modal interpretation (see Lombardi and Dieks (2014) and references therein). In the present context, it is sufficient to focus on the consistency of the property assignment in terms of *L* and *R*.

The states obtained by (anti)symmetrizing product states can thus be understood as representing particles possessing their own independent quantum properties. In this sense they are *not* entangled, in spite of the fact that the total state is not factorizable. This may seem strange since the singlet state, the paradigm example of entanglement, appears to be precisely such an antisymmetrized product. However, the complete state in Bell-type experiments is of the form (2) rather than just . This complete state is *not* the result of antisymmetrizing a product state, even though its spin part (the singlet) is. This property of the total state is responsible for the non-factorizability of joint probabilities for measurement outcomes on the two wings of the Bell experiment, and consequently for violations of the Bell inequality and for non-locality. By contrast, had we begun with a factorizable state like , we would after antisymmetrizing have arrived at

 (6)

The probabilities in this state, for spin measurements on the two wings of the experiment, *do* factorize. This means that there will be no violations of Bell inequalities and no no-go results for local models. In fact, the quantum formalism itself, with the above-mentioned provisos about the meaning of labels, immediately suggests a local model: the state (6) describes a situation in which there are two particles, one left and one right, with the particle on the left-hand side having its spin up and the right-hand particle having its spin down.

This brings us to the essential difference between the states (2) and (6). In (6) there is a strict correlation between spatial and spin properties, which is absent in (2). In terms of the particle concept that we have discussed above, we can say that (6) makes it possible to think of two particles labeled by  and , respectively (note that these differ from the standard labels 1 and 2). In (2) there is no such correlation, and this means that we cannot define physically meaningful particle labels that stand for co-instantiated sets of properties. In other words, (6) lends itself for a particle interpretation but (2) does not.

**4. EPR without particles**

The upshot of the discussion so far is that although a focus on particle labels makes it initially appear that in quantum mechanics there is no place at all for localized particles of the same kind, reflection on the meaning of the labels redresses the balance. Proper attention to what we *mean* by “particles” makes it clear that we *can* have situations of individual, even though “identical” particles, each with its own distinctive properties. We can therefore resist the idea that quantum particles are essentially non-local entities: it certainly is possible to have two identical quantum particles, one localized at  and the other at .

However, further reflection shows that exactly those situations in which we can define quantum particles by specifying sets of individuating properties, are situations that are uninteresting from a quantum information point of view. For example, the state (6) allows the definition and individuation of two particles, each characterized by its own position and spin. It is a state that respects the traditional particle concept, by tying properties together in definite bundles; but it does so at the cost of losing its quantum interest. Indeed, as we have seen, joint two-particle probabilities factorize in (6), so this state satisfies all locality requirements and will not violate any Bell inequalities. It will consequently not be suited for informational purposes that transcend classical possibilities.

By contrast, the state (2) defies attempts at defining particles by complete sets of properties: there is no correlation, in this state, between positions and spins. The very concept of a particle becomes moot in this case. Although it may still seem natural to treat *L* and *R* as a kind of “particle labels” here, since we can only measure something at these positions, these labels do not correlate to other properties than position itself. This strange situation is the one that is vital from a quantum information viewpoint, but at the same time it is a situation in which the classical particle concept breaks down.

It might be felt that this is just what was to be expected, since the violation of Bell inequalities is commonly acknowledged to demonstrate that no local account can be given of what happens in an EPR experiment −we could have anticipated from this that we should not speak of one particle at position *L* and another at *R*, each with its own properties. But this is not the moral that is usually drawn. The standard story is instead that there *are* two localized particles, but that a measurement on particle 1 instantaneously *changes* some feature of particle 2. This story presupposes the traditional concept of a particle with its own individual properties. Our analysis goes against this standard story by arguing that the traditional conceptual framework is inadequate.

Looking at (2), we see that in this state there is no correlation between spins and positions; the total state is the *product* of its spatial and its spin parts. So even if *L* is taken as a label (instead of 1 or 2), we are not allowed to refer to the “spin of the particle on the left”. Instead, there is a double-spin state that combines equally with both *L* and *R*. It might seem that on balance this returns us to a picture of the kind we have alluded to before, in which particles are non-local (present both at *L* and *R*) from the very start. But it would be inaccurate to make no distinction between what we are arguing and proposing here and that earlier suggestion of non-locality.

What we briefly considered earlier was the possibility that there are two particles but that both are non-local entities in the sense that they are present at *L* *and* *R*. But what we argue for here is the *inadequacy* of the particle concept in typical quantum cases. According to this new viewpoint, there are only two positions in space (*L* and *R*) where something can be measured; and in both places the *very same* two-spin state is “visible”. This spin state is not defined as a function on spatial points, so it is not a physical *field* (a field, e.g. an electric or magnetic field, is an assignment of field strengths to spatial points; the values of the field are therefore correlated to positions, as is clear from the notation ). By contrast, the spin state lives in its own Hilbert space, independently of the spatial part of the wave function. In other words, we should not think of spin amplitudes as being non-zero at *L* and *R* and zero at other positions, in the same way as in the case of a classical electric field that is only non-vanishing at *L* and *R*. The spins by themselves do not constitute a spatial entity at all.

Classical fields could have the same *numerical values* for their field strengths at *L* and *R*, but these field strengths would obviously not be identical in the metaphysical sense (according to which two identical things are one and the same). The proposal for the quantum case that we are considering is very different: at both *L* and *R* we can make contact with the *identically same* spin state −it is as if both *L* and *R* are windows through which we are able to look at exactly the same scene (which itself is not spatial). As it turns out, this picture leads to an explanation of the EPR-Bohm experiment that may be called local, even though this explanation is essentially non-classical and dispenses with the particle picture.

**5. Local quantum explanations**

Let us look at the details of what happens if an EPR-Bohm experiment is performed in the state (2). What usually is taken to be the phenomenon in need of explanation is that a spin measurement on the left wing of the experiment produces a definite spin value on the right wing, while there was no such definite value before and in spite of the fact that no signal propagates from one side of the experiment to the other. As we have argued, this way of formulating the problem presupposes that there is a correlation between spin values and positions −that we have particles at the two wings of the experiment that possess their own spin characteristics (non-definiteness of the spin value being such a spin characteristic), which are subsequently changed by the measurement. Our rejection of this presupposition thus dissolves the problem as originally posed. But we should of course investigate the explanatory resources of our own alternative account to see whether this change of perspective is of any help.

Let us suppose that a spin measurement in the z-direction is made by a device located in the region *L*, and that the result is “up”. Let us first consider what happens if we make use of the projection postulate (collapse of the wavefunction); this is the treatment that in the usual approaches most strongly suggests that a superluminal influence propagates between the wings of the experiment. Application of the projection postulate within our conceptual framework leads to the following account. The spin-measuring part of the device is able to make contact with the singlet state (the position-independent double spin state) via the “space window” located at *L*. The outcome “up” gets registered by the device at *L*, and as a result the singlet state collapses to the spin-down state . *This very same state* is subsequently observed from the “window” *R* when a spin measurement is undertaken at that position. There is no need for any signal between *L* and *R*, since the spin state is not spatial. The spin state is identically the same regardless of whether seen from *L* or *R*.

Formally, the treatment with the projection postulate goes as follows. The projection operator that represents the spin measurement on the left wing is

 (7)

where account has been taken of the fact that the spin measurement is at the same time a position measurement, by a localized device. The projectors representing the spin measurement are therefore correlated to position. The result of making *P* work on  of Eq.(2) is

 (8)

which is a properly antisymmetrized product state. As we have seen in Section 3, such product states represent particles in the usual sense, for which spatial properties and spin properties are bound together in an identifying set of properties. The explanatory story remains essentially the same as just sketched: the physical interaction localized at *L* transforms the singlet state to . Before the measurement, the spin state was not a spatial entity (it was not correlated to the spatial wavefunction), but the local measurement at *L* correlates “up spin” to . The “down spin” becomes consequently correlated to , not because of a signal going from *L* to *R*, but because the event of establishing the latter correlation is identical to the measurement at *L*. In this sense the explanation is *local*.

The effectively same story can also be told without collapses, employing only unitary time evolution. This is the account that should be used in order to give a local explanation of EPR-type experiments within non-collapse interpretations of quantum mechanics. In the familiar schematic von Neumann rendering, in which a measuring device possesses a neutral state  and post-measurement states  and , we obtain after a measurement interaction in which only local interactions are involved:

 (9)

So, relative to the device state  representing the outcome “up” we get a properly antisymmetrized product state according to which there is a left-hand particle with spin up and a right-hand particle with spin down. This does not require any non-local interactions precisely because the spin part of the original state , i.e. the singlet state, is uncorrelated to position. The local spin measurement at *L* in effect (if we think of the antisymmetrized product state as an effective product) selects from this singlet the part , and this is possible via an interaction that is restricted to the region *L* because the whole singlet is available there.

The mechanisms that lie at the basis of quantum information transfer can thus be understood in a local way, provided that we realize that the classical notion of a particle, or a field of properties defined on points in space, is not generally applicable in quantum theory.

**6. From classical to quantum**

In Section 5 we considered examples of transitions from a general quantum picture, in which there are no sets of interconnected properties defining particles (or sets of properties tied to a spatial point), to a classical picture in which there *are* such particle or field properties. The existence of these transitions shows that quantum theory is able to accommodate classical particles and fields as special cases, as we would expect from a theory that is more general than classical physics. But the opposite transition should also be possible: starting from a classical situation of localized systems, with their own properties, it should be feasible to produce typical quantum situations. If that were not the case, it would be impossible to design laboratory experiments in which quantum teleportation and other entanglement assisted quantum information transfer processes manifest themselves.

The simplest example of such a transition from classical to quantum is the paradigmatic Bell case, in which a localized system, initially with its own individuating properties (e.g., total spin ), decays into two parts that do not behave as particles, in the sense that they do not possess individuating properties (the definite-valued spin state  splits into two parts whose combined spins are described by the singlet state, for example). Suppose that the two spatial wave packets subsequently travel far apart due to free propagation. The spin part does not change by this since it is not spatial and unrelated to the evolution responsible for the spatial traveling. This finally leads to the EPR-Bohm state, which is not a particle state as we have seen, even though it has two spatial parts at a distance. The transition from classical to quantum took place locally here, in the decaying process, and was unaffected by the subsequent traveling of the spatial parts.

A more sophisticated experiment with entangled electrons has recently been discussed and implemented by Hensen *et al.* (2015); compare also the alternative experiment with photons by Giustina, Versteegh *et al*. (2015). The results of these experiments are important for various reasons, not the least of which is that they hold prospects for applications in quantum information protocols. In the Hensen experiment, remote electron spin entanglement is created without traveling electrons (following a proposal by Barrett and Kok (2005)). At the start of this experiment two localized electrons, both in definite spin states and therefore comparable to classical particles, are prepared at stable positions *A* and *B*, at a large distance (1.3 km) from each other. The two initial spin states are the same, namely , so that the two-electron state at this stage can be represented by

 (10)

Note that it would make no sense to say that in this state one particle finds itself at *A* and the other at *B* if we assumed that the labels 1 and 2 refer to particles. As we observed before, the standard terminology in experimental physics presupposes an analysis of the individuality of particles as we have given in section 3.

In the next step of the process the electrons at the two positions are excited by a short laser pulse, after which spontaneous photon emission entangles the local electrons spins with the photon field in such a way that the spins together with the field are described at both wings by the state , in which 1 and 0 denote the presence and absence of a photon, respectively. The total state is accordingly transformed from (10) into

 (11)

To simplify the notation we have indexed the photon/spin states with the associated electron spin labels.

Now, if a one-photon state is detected at a third position, possibly very distant from *A* and *B*, without any discrimination with respect to its provenance from either *A* or *B*, the state (11) effectively collapses to its part containing one photon ( or ):

 (12)

This is a maximally entangled spin state, comparable to the EPR-Bohm state of Eq.(2).

The actual procedure followed in the experiment by Hensen *et al.* (2015) is more sophisticated and involves a further step to eliminate the “detection loophole”, which in this case is the possibility that only one photon is detected, due to incomplete detector efficiency, whereas in reality two were emitted. This additional step guarantees that detection of a photon “heralds”, with certainty, the existence of an entangled pair of electron spins, associated with the localized spatial electron wave functions at *A* and *B*.

In this experimental procedure the initial state was classical in the sense that it described two well-localized electrons, each with a definite spin. This familiar picture became invalid when the electrons interacted with the electromagnetic field: through purely local interactions (resulting in spontaneous photon emissions) the electron spins were entangled with the photon field. After this entanglement the electron positions *A* and *B* were no longer associated with pure spin states, so that the classical particle picture ceased to be applicable: although electron position measurements could still only have success at *A* and *B*, it was no longer true that electrons with complete sets of electron properties were present at these positions. The (again purely local) measurement on the photon field, which may have taken place at a large distance from both *A* and *B*, finally fully decoupled the spins from the positions *A* and *B* so that the total state (12) became a product of spatial and spin parts.

**7. Conclusion**

The classical theories of physics, i.e. non-quantum theories up to and including general relativity, are all *space-time theories* in the sense that they define physical systems through sets of properties that are functions on a manifold of space-time points. This theoretical feature immediately implies a *basic principle of locality* that is satisfied, namely that physical systems and their histories can be completely described by the specification of all the local states of affairs in their existence. A paradigmatic example of such a classical system is a particle, in particular a point-particle: all its properties, including its position, are bound together in one package, which defines the particle and its state. This first locality principle has to be distinguished from a second one, namely that signals cannot travel faster than light. When we combine the two principles we arrive at the result that a particle at *R* cannot undergo any changes due to what happens at *L*, unless there is enough time for a causal signal to propagate from *L* to *R*.

The structure of the formalism of quantum mechanics suggests an entirely different picture. At first sight, it may seem that in the case of particles of the same kind (“identical particles”) there is no scope for the concept of a localized particle at all, because of the (anti)symmetrization rules: the particle labels are evenly distributed over all positions where the total many-particles wavefunction does not vanish. But this first obstacle for a particle interpretation can be quickly removed, by individuating the particles via their states instead of via the labels that occur in the formalism. However, this manoeuvre has only limited success, because it only leads to an analogue of the classical particle picture in quite specific circumstances. In the general quantum situation “particle properties” do not combine together to form individual particle states. In particular, it is possible to have localization in individual narrow regions in the spatial part of the total state, without complete packages of particle properties that are correlated to these individual regions. The basic locality principle just mentioned (“locality in the first sense”) therefore fails in quantum mechanics.

This does not entail that the second locality principle that we mentioned, namely the impossibility of superluminal signaling, fails too. This second principle is respected in quantum mechanics, but this is only relevant for the explanation of correlations between measurement results if we are dealing with a space-time theory in which properties are local in the sense of our basic notion of locality. This means that it is true that quantum features that are bound to position cannot be changed except by signals that have enough time to reach them, just as in classical physics. But according to the quantum formalism there is also the possibility of physical properties that are independent of space, like the spins of the singlet state. When such space-independent properties are modified by some intervention the effect of the modification is valid for the total system at once. Such a change suggests an action at a distance when interpreted as a change in *local* features; but our point is precisely that in these cases we should *not* think in terms of local properties, not even when the spatial part of the total wavefunction is restricted to well-defined narrow regions.

The new possibilities of quantum information transfer depend on this lack of *basic locality* in quantum mechanics; they do not conflict with locality in the second sense and do not conflict with relativity theory. Awareness of this double role of locality and the precise way in which quantum mechanics is non-local appears to provide a framework for understanding the effectiveness of quantum information processes.

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