Quantum Measurement

Subjects: Others
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The Special State Theory of Quantum Measurement

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There is only pure unitary time evolution in the special state theory of quantum measurement. In this, it resembles the Many Worlds Interpretation. Nevertheless, there is only one world. This is accomplished by means of “special states” and significant modification the arrow of time. Experimental tests are described, although they have yet to be done.

Keywords: special states ; arrow of time ; foundational experiments ; statistics

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The interpretation of quantum measurement has been a puzzle for a century. For most applications the mystery can be ignored. But when one deals with the macroscopic and microscopic simultaneously, in particular when a “*measurement*” takes place, there still rage disputes.

Proposing a solution to such a long-standing puzzle requires, as Bohr said,[2] “a sufficiently crazy theory.” What is presented below satisfies Bohr’s criterion. In particular, my ideas depend on two radical ideas: the existence of special states and modification of the arrow of time. I myself have trouble believing these ideas, but what motivates me is that other theories of the measurement process are even more difficult to accept.

In the special state theory of quantum measurement *measurement* is just unitary time evolution. Since the early days of quantum mechanics experimental technique has blurred the line between macro and micro, and there has never been any indication of any dynamics but unitary time evolution. The special state theory thus shares a common view with the Many Worlds Interpretation. The difference is that I have only one world. This requires significant modification of statistical mechanics, as well as the existence of particular quantum states, states that I call *special*. This is *not an interpretation* and is physically different from other theories. In particular, there are experimental tests and some are described here.

Special states in quantum mechanics. To see how *special states* operate, consider an example.

Suppose there is a single 2-state system in contact with a heat bath of bosons [12] . Initially the system is in its excited state (spin-up) and the boson coupling induces the decay. The idea is that when the the two qubit state is in its lower level it has measured the presence of the boson system. This model of measurement misses much of the real world, in particular “ registering” the measurement, making sure the system does not return to its excited state. These features are assumed to result from parts of the system that I do not model. I use the spin boson model with a single boson, and the Hamiltonian can be taken to be

$$H = \frac{\varepsilon}{2}(1+\sigma_z) + \omega a^\dagger a + \beta\sigma_x(a^\dagger + a)\,. \qquad (1)$$

The Pauli spin matrices are the operators for the 2-state (spin) system, a and $a^\dagger$ are the boson operators and ε, β and ω are parameters.

As indicated, the single spin starts (say) as $\psi_{\text{up}} \equiv \binom{1}{0}$, and the coupling to the bath tends to flip it over. The *special states* are particular initial conditions of the bath such that the *microscopic* final state of the spin is either **up** or **down**, i.e., $\binom{e^{i\phi}}{0}$ or $\binom{0}{e^{i\phi}}$ for real ϕ. “Final” in this case refers to a specific time, namely that time at which a large, irreversible system notes the state of the spin. Such “special” states would be rare, since for most initial states after a time the microscopic state would be partly up, partly down, something of the form $\binom{\alpha}{\beta}$ for nonzero α and β.

Let’s be specific. The system has Hamiltonian (1) with ϵ= 0.5, ω = 0.1, and β = 0.6. The spin starts in $\psi_{\text{spin}}(0) = \psi_{\text{up}} \equiv \binom{1}{0}$, with the oscillator state unspecified. The finiteness of the computer requires that the oscillator be cut off, and we only considered 250 states. (This led to an error in [a,$a^\dagger$] only in the 250$^{\text{th}}$ diagonal term. To reduce this effect only states with small probability in the highest levels were considered.) For these parameters at time-0.15 (which is the *specific time*, the time at which a large irreversible "observer"[3] looks at the spin) there is about a 50% probability of decay, namely if one would trace out over the oscillator states the density matrix for the spin would be half-half.

But what concerns us is whether there are states that have entirely decayed, or not decayed at all. The way to learn this [13] is to define a projection operator on the spin: $P \equiv (\psi_{up}\psi_{up}^{\dagger}) \otimes \mathbb{1}_{\text{boson bath}}$. Using this operator, the probability of being all **up** at time *t* is

$$\Pr(\mathbf{up}) = \langle\psi_{\text{up}} \otimes \psi_{\text{bath}}|U^{\dagger}PU|\psi_{\text{up}} \otimes \psi_{\text{bath}}\rangle = \langle\psi_{\text{up}} \otimes \psi_{\text{bath}}|PU^{\dagger}PPUP|\psi_{\text{up}} \otimes \psi_{\text{bath}}\rangle\,, \qquad (2)$$

with $U \equiv \exp(-iHt/\hbar)$ and where I have used *PP* = *P*. Defining *A* ≡ *PUP* and using $P^{\dagger} = P$,we have $\Pr(\mathbf{up}) = \langle\psi_{\text{up}} \otimes \psi_{\text{bath}}|A^{\dagger}A|\psi_{\text{up}} \otimes \psi_{\text{bath}}\rangle$. Defining $B \equiv A^{\dagger}A$ (a function of *t*), it follows that the issue of whether any initial state (of the bath) can lead to a measurement of **up**, using purely unitary time evolution becomes the issue of whether *B* has eigenvalues equal to one. Similarly the issue of whether there any fully decayed states requires of B that it have eigenvalue zero. (B's eigenvalues are real and in the interval [0,1].)

Indeed, there are such states. One probability distribution for each case is shown in Fig. 1. The figure only indicates magnitudes; the relative phases are fixed, but not shown. Thus the state that at time *t* = 0.15 is *entirely* in the up position had an initial bath state of the form $\sum c_n exp(i\emptyset_n|n\rangle$, with the values of $\{|c_n|^2\}$ shown in Fig. 1a (and the particular values of $\{\emptyset_n\}$, not shown). At time-0.15 this state is undecayed *despite* the fact that a random or average state would be a superposition of both decayed and undecayed states, with the trace (over bath states) giving about 1/2 for each eventuality. Similarly, the state whose probabilities are shown in Fig. 1b has decayed almost completely at this time. (There is a slight residual probability both for decay and non-decay. In the illustration Pr(decay) for the "non-decay" state is about 6.1×10^{-4}, while Pr(non-decay) of the "decay" state ≈ 8.1×10^{-5}. This will be discussed below.)

What good are these states? They constitute the main idea for avoiding *many* worlds while holding on to unitary time evolution.

Suppose a (Schrodinger) cat is placed in a chamber with the usual vial of poison. Let the dispersal of the poison be governed by the spin state just discussed. The spin-cum-bath system is also in the chamber and the entire setup isolated. Its isolation ceases at time-0.15 and their reversible "registration" is due to the observer, who looks to see if the cat is dead or alive. In the usual theory there is significant probability for both options. However, there is no problem if the initial state of the bath is one of the "special" states I have been describing. For non-decay initial conditions there is a living cat, for decay initial conditions it is dead. This is accomplished with no black magic; it is the result of unitary time evolution from "special" initial conditions.

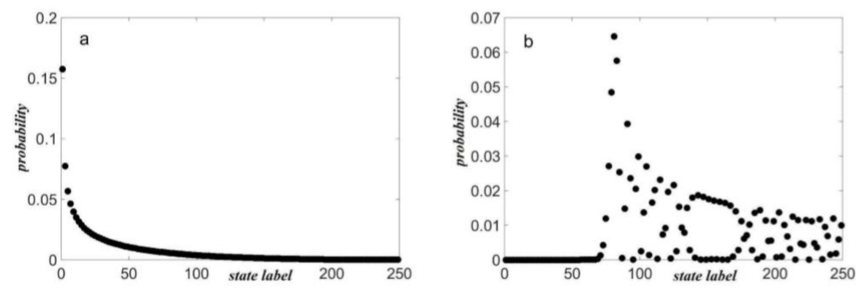

Figure 1: Special *time*-0 oscillator states. Figure (**a**) shows the (initial) probability of excitation of oscillator states that contribute to the non-decay state. Only shown are even states, since there are no odd states. Phases are not shown, but are also fixed by the non-decay condition. In Figure (**b**) are shown the probabilities for the state that decays; in this case only even oscillator states are shown (the rest are zero). As in image a, the phases, though not shown are crucial to the "special" nature of the state.

That is the main idea of the special state theory: no macroscopic superpositions because of particular initial conditions. There is also no entanglement. At time-0.15 the spin state is wholly in one state or the other and a trace would leave the oscillator state unchanged.

The next question is, why should Nature arrange to have a "special" state as the initial conditions for every situation where a potential split into many worlds occurs? Moreover, "'final" states for one measurement are (special) initial conditions for the next. This implies both an extreme form of determinism and a sufficient abundance of special states. I don’t have justification for these assumptions, except to say that this is the conclusion I am driven to by insisting that no magic dynamics occurs in the measurement process and that there is only one world. What I can offer though is perspective. How strange is it for there to be particular, non-random, initial conditions? For this I turn to the next topic.

Thermodynamics and the arrow of time. The usual arrow of time is equivalent to using random initial conditions [13] . This has never been verified experimentally; its main virtue is that answers computed based on this agree with experiment —not to be sneered at, but not a proof.

I use a model of mixing dynamics to show that random initial conditions are unnecessary. The model is the cat map [6] , a transformation of the unit square (coordinates x,y) into itself

$$\left.\begin{aligned} x' &\equiv x+y \\ y' &\equiv x+2y \end{aligned}\right\} \bmod 1\,, \qquad \text{or } \begin{pmatrix} x' \\ y' \end{pmatrix} \equiv M \begin{pmatrix} x \\ y \end{pmatrix} \bmod 1 \text{ with } M = \begin{pmatrix} 1 & 1 \\ 1 & 2 \end{pmatrix} . \qquad (3)$$

Imagine an ideal gas of N such points in the unit square, each moving in discrete time under the cat map. In Fig. 2 I show what happens to a collection of $N = 500$ points initially satisfying $0.5 \le x \le 0.6$ and $0.5 \le y \le 0.6$. This mechanical system is headed for chaos. To be quantitative about this “chaos” define an information entropy using a coarse graining. As grains take the 100 1/10 by 1/10 squares contained in the unit square, and only count the number of points in each grain. Thus the observer can only determine which grain a point is in, not the point’s exact coordinates. The entropy is

$$S = -\sum p_k \log p_k\ , \qquad p_k \equiv n_k/N\ , \qquad (4)$$

where k runs over the coarse grains and nk is the number of points in grain-k. The behavior of the entropy for the points in Fig. 2 is shown in Fig. 3.

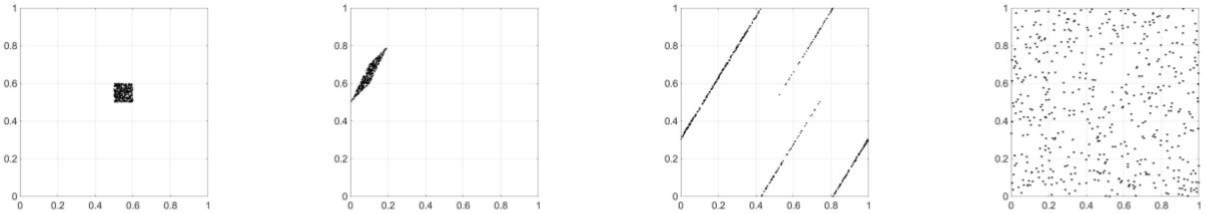

Figure 2: 500 points are started in the square (x,y) ∈ [.5,.6]⊗[.5,.6], but are otherwise random. Under the cat map, Eq. (3), they become the parallelogram in the next figure. They have stretched in the direction of the eigenvector with larger eigenvalue of M (Eq. (3)) and shrunk along the other eigenvector. (The dynamics satisfies Liouville’s theorem because det $M = 1$.) On the next time step (the next figure) the mod 1 action coupled with the stretching has begun to pull the points apart and by time 7 (the last figure) nothing is recognizable.

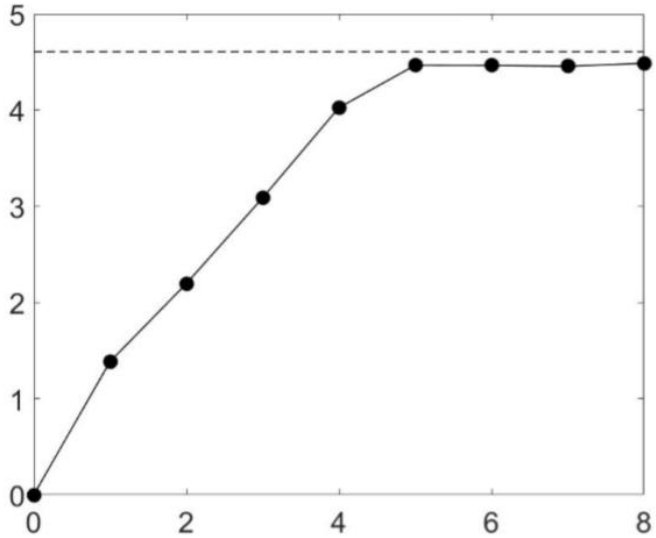

Figure 3: Entropy, as defined in Eq. (4), as a function of time for the simulation of Fig. 2. For this grain size, equilibration sets in at about time-5. The dashed figure is the maximum entropy for this coarse graining (log100). It is not attained because the number of points is finite.

This is a manifestation of the normal arrow of time: the initial points were selected randomly within a particular grain and the entropy increases monotonically until close to equilibrium, after which it fluctuates in a predictable way.

But let's consider another simulation. In this case the initial points were not selected randomly, but for the first time steps it will look that way. The simulation runs for 18 time steps away from the initial square and is shown in Fig. 4. The sequence of images should be read left-to-right and row-by-row. There are 4000 points and most time steps are illustrated. *Every point in this simulation evolves by pure cat map dynamics*. So how did I get it to give me these strange images? It was easy. I randomly occupied the first little square ($0.6 \le x \le 0.8$ and $0.4 \le y \le 0.6$) at time-0 with about 30×4000 points. Then I imposed two conditions on the points. First, at time-18 they needed to occupy a different little square of the same size as the original. This cut the number of acceptable points by a factor 25. Then I also required that at time-9 the points arranged themselves in the figure of a cat. This cut things down a bit more, so I was left with 4000 points that satisfied all 3 conditions. This solved a 3-time boundary value problem. It was trivial in this case since the points were non-interacting.

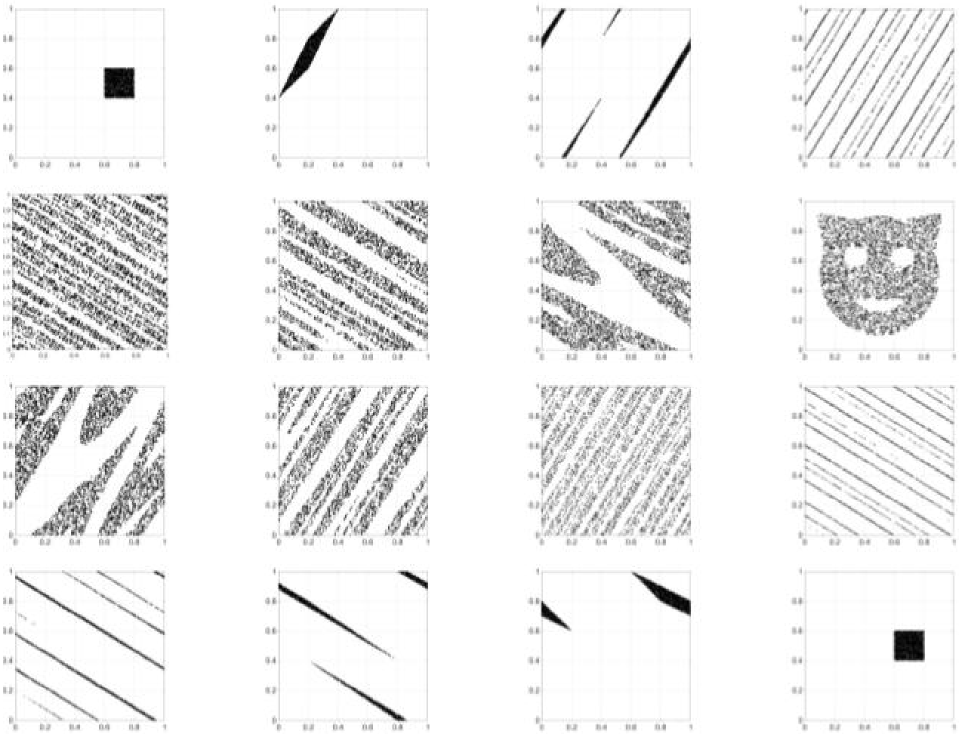

Figure 4: The times are as follows: row-1: [0, 1, 2, 4]; row-2: [6, 7, 8, 9]; row-3: [10, 11, 12, 14]; row-4: [15, 16, 17, 18]. These points all evolve under the cat map and satisfy boundary conditions at times 0, 9 and 18. The size of the position marker (a point) varies with the image, for better visibility.

The message of this demonstration is most poignant when phrased in terms of entropy. Grain sizes are now 1/5 by 1/5. For each of the configurations shown the entropy was computed and is graphed as the circles in Fig. 5. A second curve, with star markers, is also shown in that figure. It is the entropy for 4000 points initially starting in the same coarse grain, but having no additional requirements. (Like Fig. 3, it corresponds to the usual purely random initial conditions.) Compare these two curves. The point is that <u>for times prior to about 7 you **cannot** tell the difference</u>.

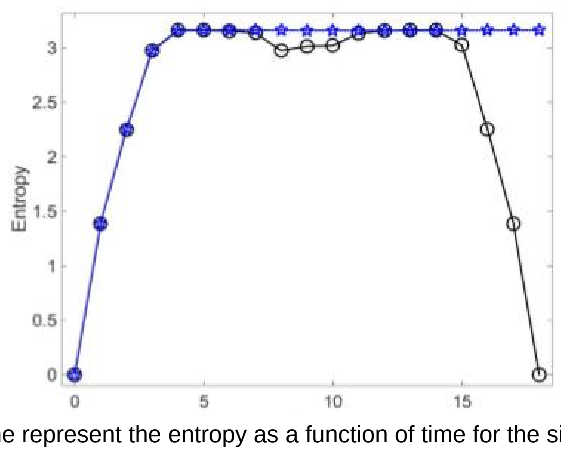

Figure 5: The circles and solid line represent the entropy as a function of time for the simulation shown in Fig. 4. Note that it drops a bit at time-9 and then at time-18 plummets back to 0, since all points are again in a single grain. The other curve, marked by stars, is the entropy as a function of time for 4000 points having the same initial grain, but with no other constraints.

My message is simple: there may be future constraints, but you would not know about it. The arrow of time does not (necessarily) point as fixedly as one might have supposed.

Another way to say this is to observe that the points in the simulation that yielded Fig. 4 were not random but had a *cryptic* constraint, a constraint that was difficult to discern, but which nevertheless plays in important role as the dynamics unfolds.

My objective is to justify the idea that there could be cryptic constraints in the world. Not every imaginable state occurs in Nature; only those which, in my terminology, are special. This is a severe restriction, but I have demonstrated that the restriction may be invisible. The kind of restriction that I find most palatable is a two-time boundary condition. This contradicts the usual arrow of time, but its effect may not be noticed except close to the boundaries.

What kind of two-time boundary condition could select special states? First consider initial conditions. As Wald [19] has pointed out, in the early universe the entropy was low, for unknown reasons. I go a step further: assume the von Neumann entropy was also low; there was little or no entanglement. Next, imagine that our entire cosmology is roughly time symmetric. This idea is not popular today due to the apparent accelerated expansion. But that phenomenon is poorly understood and there have been suggestions of a periodic cosmology despite the acceleration (a sample is [10, 18]). One additional component enters, the connection first suggested by Gold [9] , relating the arrow of time to an expanding universe. One then expects that under contraction the arrow will be reversed. (This my assertion; Gold has denied saying this.) Recalling the connection between boundary conditions and the arrow of time, one can now enunciate a possible boundary condition that would demand special states: no entanglement at the beginning, no entanglement at the end.[4]

To see why this invokes special states, consider a Schrödinger cat. At the end of the experiment—and again demanding only unitary time evolution—in the MWI there is a portion of the wave function with a dead cat and a portion of the wave function having support on a macroscopic state recognized as a living cat. The dead one is buried, the living one continues to hunt mice. But how can these portions of the wave function be recombined coherently, as they would need to be if there is a no-entanglement demand in our future. It would take tremendous coordination to accomplish this coherently. Having a special state is also an unlikely way to avoid entanglement, but it is *less* unlikely than recombining a superposition of macroscopically distinct states.

With the two-time boundary rationale one can approach the issue of the small amount of leftover wave function. As indicated, the special state for decay has a small but non-zero probability of non-decay (in the examples it is roughly 10^{-4}). In the context of a boundary value problem one does not need perfection. The measure of possible error in “specializing” is given by the tolerance of that boundary value problem. I also expect that even this level of error is much larger than what occurs in nature and represents a limitation of my computer.

Recovering the Born probabilities. If every experiment involves interaction with apparatus and special states, why is that probabilities can be calculated using only the wave function of the system being studied? This is our next subject.

In the usual interpretations of quantum mechanics you don’t pay attention to the apparatus: to calculate probabilities you only need the wave function of the system being measured. By contrast, the special state theory must also construct the wave function of the apparatus. It therefore seems necessary to make the following claim: the abundance of special states of the apparatus reflects the wave function of the system being measured. For example, suppose you have two outcomes to an experiment, with respective amplitudes α and β ($|\alpha|^2 + |\beta|^2 = 1$). Then the number of special states (which requires looking at a larger space) for each outcome will correspond to the same relative probabilities, $|\alpha|^2$ and $|\beta|^2$. Although the space of special states is small compared to that of all states, nevertheless the vector space for each outcome is still comprised of many dimensions and the relative dimensions reflect the usual probabilities. In this way “probability” in quantum mechanics is like probability in classical mechanics (using the correspondence of volume in phase space with dimension in Hilbert space). The outcome is determined (by unitary time evolution), but since you don’t know which special state is involved, you must fall back on probabilities. This is not a hidden variable theory. The ignorance of which special state is involved reflects the macroscopic nature of apparatus. In principle, if one knew the *microscopic* state, one would know the outcome with certainty.

I will come back to this claim later. For now I will establish a weaker result. The actual special states will not be identified, but their possible statistics will be studied.

Consider the simplest quantum measurement problem, determining the state of a two-level system. For convenience we study a Stern-Gerlach apparatus measuring the spin along the z-axis. Let the initial spin state be

$$u_\theta = e^{i\theta\sigma_x/2}\begin{pmatrix}1\\0\end{pmatrix}. \qquad (5)$$

(In this basis, up means an eigenstate of the z-component of spin with a magnitude 1 entry in the upper portion. Correspondingly for **down**.) There are of course other degrees of freedom in this problem. Foremost is the position of the spin (riding on, say, a **K** atom); then there is the macroscopic magnet, macroscopic screen, and more. But we are now working backwards and will assume that lurking in various corners of this enormous Hilbert space are special states for the outcomes, **up** or **down**. It is the statistics of these states that we study. In particular, we will consider the abundance of different special states for different angles, θ, of the spin orientation. At first it will seem that this abundance issue is impossible to satisfy, but then it will be seen to impose a particular structure on the space of all spacial states.

The first step in any assessment of the availability of special states is to seek the least unlikely way in which they could occur. Of course they are rare, but there are different levels of rarity and we seek the kind of special state that imposes the fewest conditions on the wave function. For an Stern-Gerlach experiment the atom must ultimately be absorbed in one or another region of the target area, but not in between (or elsewhere). This can be accomplished by combining the output wave functions after they've left the magnet, or by rotating the spin so there is only one wave packet that emerges from the magnet (I assume familiarity with the Stern-Gerlach setup). Rough estimates (see [17]) suggest that recombining the spins is far less achievable than rotating the spin before it even enters the magnet. So the problem becomes one of rotating the spin from its initial value (θ) to be **up** or **down**.

What one might take to be a stray magnetic field provides just the right force to bring the spin to 0 or π, prior to its entry into the inhomogeneous field. As usual, the function of the inhomogeneous field is to make the position coordinate dependent on the spin with which the atom entered that field.

Let us call this apparently (but not truly) random change in θ a "kick." If you imagine doing several experiments with different initial θ's, the number of kicks to one or the other special value also varies. Let $f(\theta)$ be the probability of obtaining a kick of size θ. For any one experiment only kicks of size $-\theta + n\pi$ (with n an integer) enter, but it is reasonable to assume that as θ varies there is some well-defined distribution, manifested in our situation (due to our ignorance) as a probability. Thus to get spin up one must have a kick of size $2n\pi-\theta$ for positive or negative (integer) n's. Similarly, to get spin down one requires a kick of size $(2n + 1)\pi - \theta$. Thus the probability of spin up is $g(\theta) = \sum_{n=-\infty}^{n=\infty} f(-\theta + 2n\pi)$. Similarly for spin down one adds π to eachsummand in the argument of f. On the other hand, standard quantum mechanics, i.e., the Born rules, dictate that the ratio of down to up is $\tan^2(\theta/2)$. Therefore our requirement on f (hence on g) is

$$\tan^2\left(\frac{\theta}{2}\right) = \frac{g(\theta+\pi)}{g(\theta)}\,, \quad \text{with} \quad g(\theta) \equiv \sum_{n=-\infty}^{n=\infty} f(\theta + 2n\pi)\,. \qquad (6)$$

(In Eq. (6), in the definition of g, use has been made of f's $\theta \to -\theta$ symmetry as well as the fact that n is a dummy variable.) The function $f_0(\theta) = 1/\theta^2$ is an explicit solution to Eq. (6), but is not normalizable, as a probability should be. (There is no normalizable solution to Eq. (6).) However, for θ close to 0 it is possible to cut off the function $1/\theta^2$. A convenient cutoff makes use of the Cauchy distribution

$$C_a(\phi) = \frac{a/\pi}{a^2+\phi^2}\,, \qquad (7)$$

which for small enough a changes Eq. (6) very little. The deviations from standard probabilities are largest for $\theta \sim 0$ and are of order a^2; since a is unknown, one can only bound it, since as far as I know deviations from the standard results have not been found. The distribution C_a does not have a second moment. But I emphasize that for the purposes of the experimental tests described below and in [17] , the Cauchy distribution is not required.

Experiments. In [13] experimental tests were proposed, but they were more for the purpose of demonstrating that the ideas were different physics, not another interpretation. Another possibility came up in dealing with the time it took for a quantum system to change states [15] , but there too the practicality was not clear.

More recently experiments were proposed that should be feasible. They are challenging and we have not yet convinced anyone to carry them out, but I believe they are doable. Here I review an experiment involving a modified Stern-Gerlach apparatus.

Force-free rotation? *two* strong, inhomogeneous fields, one to orient along a particular direction (and half removed), and the second to measure the spins of those prepared in this way. Stern undertook such experiments, but [11] ran into trouble because unless the magnetic field between the magnets vanishes the spins follow the residual field adiabatically.

Later, in Stern's laboratory, Frisch and Segre [1] overcame this problem and they published their results jointly [8] . Frisch and Segre used work of Majorana [2] to confirm that the experiment was theoretically successful. (Majorana's analysis is analytic, but it's amusing that in the 21st century one could use a computer to get slightly sharper results (unpublished).

For simplicity assume that the two strong magnetic fields sort along perpendicular directions, so that (choosing axes) the initial wave function is the plus eigenstate of $Jy(so\theta = \pi/4$ for the θ of Eq. (5)). At the end of the experiment we get **up** or **down** eigenstates of J_z. Does this mean a force acted to rotate the spin?

In various interpretations(Many Worlds, Copenhagen) there is no force needed to go from an eigenfunction of J_y to an eigenfunction of J_z. But in the special state theory, there is!

In Many Worlds, individual observers change their perception of the value of the angular momentum, *but* there is no need for angular momentum non-conservation, since the total Hamiltonian (including the observer) commutes with **J**. The explanation is slightly different for (my understanding of) "Copenhagen." Until you actually measure J_z it has no value, since J_z does not commute with the projector for the spin state.

However, with only one world—the contention of the special state theory—there can be no change in the wave function without a proximate cause. If a quantity is changed the single observer can, if it is practically possible, determine what caused the change. If $\langle |J_z| \rangle$ (of the spin) changes its value, something else has to pick it up. This "something" can only be due to the peculiarities of the special state, what has been called a kick earlier. (Recall, the kick is not a deviation from the laws of nature, but like the cat at time-9 in the progression of Fig. 4, it follows the rules, but happens because of unusual initial conditions.)

As discussed, the least unlikely form of the kick is that a short term magnetic force acts to bring uθ to $|\uparrow\rangle$ or $|\downarrow\rangle$ (up to phase factors). The first point to investigate is how large a field is needed. Unfortunately we do not know much about the special state that does this rotating; presumably it involves quite a few degrees of freedom. However, we do know that in the end it must affect the spin through its interaction Hamiltonian which is $-\boldsymbol{\mu}\cdot\boldsymbol{B}$. What will be demanded of the field is that it causes the angle in Eq. (1) to rotate by an angle (with n an integer) in order to bring the spin entirely to an up or down state. Therefore we require that φ be at least of order unity, i.e.,

$$|\phi| = \left|\frac{-\boldsymbol{\mu}\cdot\boldsymbol{B}\Delta t}{\hbar}\right| = \mathrm{O}(1)\,, \tag{8}$$

where Δt is the time during which the special state acts. This assumes that **B** is roughly constant during its action. From the known values of $\hbar$ and $\boldsymbol{\mu}$ this requires that

$$B\Delta t \geq 10^{-11}\,\mathrm{Ts}\,. \tag{9}$$

This is a lower bound. If, as argued in [16], the φ's are Cauchy distributed then the distribution has neither a first nor a second moment. This implies that occasionally there will be large values of n (hence B) occurring, making detection easier.

Without knowing the source of the $\boldsymbol{B}$ field (in Eq. (9)) it is difficult to estimate Δt, but whatever it is, it must not be longer than the traversal time from one Stern-Gerlach setup to the next. The distance scale is centimeters (call it L and take L = 10cm) and velocities are about 500 m/s,[5], so that $\Delta t \leq 2\times10^{-4}$ s. Using Eq. (9) this implies that $\boldsymbol{B}$ must be at least 5×10^{-8} T or about 0.0005 gauss. This is about one thousandth of the earth's field, but is well within the range of measurement. Presumably this field is also time-dependent, so that by Maxwell's equations there is necessarily an electric field. From $\boldsymbol{\nabla} \times \mathbf{E} = -\partial B / \partial t$ we make the rough estimate $E \sim BL/\Delta t \sim Bv \sim 2.5\times10^{-5}$ V/m.

Theoretical issues.

• Identical particles. Phase space is reduced, since there are fewer possible states. But all particles are identical so that one "particle" can substitute for another.

• Non-zero amplitude for "wrong" outcomes. What happens in reality, without computational limitations?

• For the given example of a special state the probabilities of relative outcomes were halfhalf. If they were 70%-30%; matching the Born probabilities requires that the special states come in that ratio. In computer calculations this has not been the case. This may mean the whole story is wrong, or that the distribution of really special states has not been probed.

• The result in [14] suggest a cessation of entanglement. This is a vote in favor of special states.

Conclusions

The theory outlined above presents a coherent view of quantum measurement. At present the special state theory is seldom included in lists of quantum measurement ideas, although there is no experimental or theoretical evidence against it. My own take on this is that it advances two radical ideas. A good science fiction story only has one, and my story is not even supposed to be fiction.

The first radical idea is the existence of special states, states that lead to definite results for a measurement, rather than a superposition of macroscopically distinct states. This is a technical issue, a question that can be addressed within a well-defined context, namely quantum mechanics. To the extent that the idea has been tested numerically and analytically it seems OK.

The second idea concerns modifications of the arrow of time, determinism and even violation of the second law of thermodynamics. The fact that there are some special states also seems evident. But I do have two major concerns: are there enough special states and do they recover the Born probabilities? The suggestion of a doable experiments is heartening. What is not so heartening is that no one has yet performed them. For an experimentalist who must keep the grant money flowing there may not be sufficient motivation to divert resources to this difficult version of the Stern-Gerlach experiment or its optical analog [17] .

Note

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[2]Apparently Niels Bohr expressed this \requirement" on several occasions|and not only in connection with quantum measurement theory. See https://en.wikiquote.org/wiki/Niels Bohr (Jan. 2020 version) for variants and citations.

[3]The \observer" need not be human. The word refers to anything that irreversibly records the spin state.

[4]Of course there is entanglement every time a bound state forms. The assumption is that this is accomplished by means of special states. The electron and proton (say) denitely come together or denitely do not. As usual this can only be accomplished with the aid of other degrees of freedom.

[5]The numerical estimates are based on information in references [3, 4, 5, 7], which are not all the same. References [3, 4] are for contemporary students doing experiment who use K atoms and lower temperatures. On the other hand, [5] and [7] tell the story of the original Stern-Gerlach experiment, which involved Ag atoms.

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