J. Opt. B: Quantum Semiclass. Opt. **7** (2005) S40–S46 doi:10.1088/1464-4266/7/3/006

Quantum vacuum friction

P C W Davies

Australian Centre for Astrobiology, Macquarie University, New South Wales, 2109, Australia

Received 25 October 2004, accepted for publication 7 December 2004
Published 1 March 2005
Online at stacks.iop.org/JOptB/7/S40

Abstract
The quantum vacuum may in certain circumstances be regarded as a type of fluid medium, or aether, exhibiting energy density, pressure, stress and friction. Vacuum friction may be thought of as being responsible for the spontaneous creation of particles from the vacuum state when the system is non-stationary. Examples include the expanding universe, rotating black holes, moving mirrors, atoms passing close to surfaces, and the activities of sub-cellular biosystems. The concept of vacuum friction will be reviewed and illustrated, and some suggestions for future experiments made.

Keywords: quantum vacuum, cosmological particle creation, black holes, moving mirrors, accelerated observers

1. Casimir's legacy: the quantum vacuum rules the universe!

When Hendrik Casimir published his celebrated paper in 1948, the Casimir effect was little more than a curiosity. To be sure, the effect has subsequently been measured experimentally, but because the Casimir force is so small, it joined the Lamb shift and the anomalous magnetic moment of the electron as an elegant and crucial verification of quantum vacuum effects, but with little practical or natural significance. Today, however, physicists and cosmologists agree that the quantum vacuum holds the key to the universe. It seems very likely that the cosmos, at least as we know it, was born in a quantum event, and that the large-scale structure of the universe was fashioned by quantum vacuum processes during the first 10^{-32} s. It also seems likely that the ultimate fate of the universe will be determined by quantum vacuum energy—now dubbed 'dark energy'.

The importance of Casimir's foundational paper is that it showed how geometrical or topological constraints can affect the quantum vacuum. The significance of this for cosmology was first recognized twenty years later by Sakharov (1967, 1968), who proposed a theory of 'induced gravity' in which the modes of a quantum field in its vacuum state are changed, not from the presence of reflecting boundaries, but by spacetime curvature. The resulting shift in vacuum energy was then interpreted by Sakharov as gravitation. In this way, the static and dynamic Casimir effects would occupy a central position in physics as the ultimate source of gravitation. Most theorists, however, preferred to consider the gravitational field as fundamental and autonomous, rather than a derived product of the quantum vacuum. From this latter assumption, the subject of quantum field theory in curved spacetime was developed in the 1970s (see, for example Birrell and Davies 1982). In this theory, the gravitational field is treated as a classical background curved spacetime inhabited by quantum fields of various sorts. Einstein's gravitational field equations then become, in standard notation,

$$R_{\mu\nu} - \tfrac{1}{2}Rg_{\mu\nu} + \Lambda g_{\mu\nu} = -(8\pi G/c^2)\langle T_{\mu\nu}\rangle \qquad (1)$$

where $\langle T_{\mu\nu}\rangle$ is the expectation value of the stress–energy–momentum tensor evaluated in the quantum state appropriate to the spacetime and Λ is the cosmological constant. The definition of the vacuum state in a general curved spacetime is a subtle and complicated matter that I shall not elaborate on here. (For a full discussion see Birrell and Davies 1982.) Assuming that such a vacuum state is defined, the distorting effect of the spacetime curvature alters the stress–energy–momentum tensor of the vacuum state and produces a variety of important physical effects, the most famous of which is Hawking's black hole evaporation process (Hawking 1975). Of course, quantum field theory in curved spacetime cannot be a fundamental theory because the gravitational field itself remains un-quantized. Rather, it should be regarded as a semi-classical half-way house to a full theory of quantum gravity. In spite of this limitation, I shall restrict my treatment here to this semi-classical theory.

An early cosmological application of the semi-classical theory was to adapt Casimir's original calculation to the case of a static Einstein universe (Ford 1975). This is a spacetime in which the spatial sections are uniform 3-spheres (closed but unbounded space). Here, the closed, finite, space geometry has the effect of discretizing the modes of the field in analogy to the reflecting boundaries in the original Casimir effect.

The calculation otherwise follows Casimir's original treatment, including a cut-off at high frequencies in the mode sum. The resulting uniform renormalized vacuum energy ρ and pressure p is given for a massless scalar field by

$$\rho = \hbar c/480\pi^2 a^4 \quad (2)$$

$$p = -\tfrac{1}{3}\rho \quad (3)$$

where a is the cosmological scale factor (in effect, the radius of the universe). The electromagnetic results are similar.

The main significance of quantum vacuum energy is the fact that it mimics the cosmological constant Λ. That is, dark energy may be thought of as the energy of the vacuum. To see why this is so, ignore for the moment spacetime curvature, and consider the expectation value of the stress–energy–momentum tensor in the standard unbounded Minkowski space quantum vacuum state of the electromagnetic or massless scalar field. Then it follows from the spacetime symmetries of Minkowski space (which implies Poincaré invariance of the vacuum state) that

$$\langle T_{\mu\nu}\rangle \propto g_{\mu\nu}. \quad (4)$$

Substituting this result into equation (1) and moving the resulting term over to the left-hand side of the equation, it is obvious from the general form of equation (2) that the quantum vacuum has the same form as the cosmological constant term. As is well known, without a cut off, the vacuum energy (and pressure) is formally infinite, implying an infinite renormalization of the cosmological constant. If a cut off is inserted at the Planck energy—the natural scale for quantum gravity—the contribution of the quantum vacuum to Λ is about 120 powers of ten in excess of the observed upper limit for Λ. This staggering mismatch between theory and observation is known as 'the cosmological constant problem', and it remains one of the great outstanding challenges to fundamental physical theory.

If the only contribution to the gravitational field comes from Λ-type terms, then the gravitational field equation may be solved for the cosmological scale factor as a function of time to yield

$$a(t) \propto e^{Ht} \quad (5)$$

where H is a constant (in fact, the Hubble constant). This solution is known as de Sitter space. At first sight, one might imagine that this expanding space would result in a spectacular dynamic Casimir effect, with the quantum vacuum being exponentially stretched. However, de Sitter space has the same number of spacetime symmetries as Minkowski space, and equation (4) remains valid (Bunch and Davies 1978). Thus the quantum vacuum energy of de Sitter space remains constant even as the universe expands, and de Sitter space is therefore a self-consistent solution of equation (1).

In the now-standard inflationary scenario of the big bang theory, the very early universe is in the vacuum state of a nonlinear 'inflaton' field that generates a very large effective cosmological constant (Guth 1981, Linde 1990). This drives a period of exponential expansion with an e-folding time of about 10^{-34} s, lasting for a total duration of about 10^{-32} s, during which time the universe becomes homogenized and isotropized. This mechanism explains the large-scale uniformity of the universe observed today, and manifested especially in the cosmic microwave background radiation, which is isotropic to one part in 10^5. However, like all vacuum states, the de Sitter vacuum state of the inflaton field will suffer quantum noise. This vacuum noise will result in quantum fluctuations in the expansion rate around the average value H in equation (5). The result of this is crucial for the future evolution of the universe. Inflation ceases when the inflaton field undergoes a phase transition to a lower-energy broken-symmetry state, releasing its energy in the form of heat, which in turn creates ordinary matter. As a result of the quantum fluctuations during the inflationary phase, when inflation ceases, these fluctuations will leave their imprint in the distribution of matter in the universe (for a recent review see Langlois 2004). Specifically, there will be variations in density of this matter superimposed on the overall uniform state, and possessing a characteristic power spectrum. This prediction of inflationary theory has been well verified by the results of the COBE and WMAP satellites (Schwarz *et al* 2004). The density variations acted as the 'seeds' for the growth of large-scale structure in the universe, which eventually congealed into clusters of galaxies.

Although the end of the inflationary era was accompanied by a drop in Λ by very many orders of magnitude, there is no known fundamental reason why the post-inflation value of Λ has to be precisely zero. Recent observations suggest that the current value of Λ is of the order 10^{-3} (eV)^4, which means that the energy of the quantum vacuum is actually the dominant form of energy in the universe, outweighing ordinary matter by more than an order of magnitude (for a recent assessment see, for example Gong and Duan 2004). Because the energy density of ordinary matter falls with time, whereas the vacuum energy remains constant, then the latter will eventually come to dominate, and the universe will approach de Sitter space once again, with exponential expansion as shown in equation (5), albeit with a value of H very many orders of magnitude smaller than it was during the inflationary phase. Because de Sitter space possesses an event horizon, the end state of the universe will be one of dark emptiness. Thus nothing less than the ultimate fate of the universe rests with the energy of the quantum vacuum.

2. Vacuum viscosity and particle creation by the expanding universe

The examples discussed in the foregoing section—Minkowski, Einstein and de Sitter spaces—are all cases of spacetimes for which there exists a natural definition of the field modes and hence a natural definition of the quantum vacuum. The symmetries of the spacetime ensure that these vacuum states are stable in the case of linear fields; in particular, the electromagnetic vacuum is stable. More generally, the vacuum state will not be stable: particles will be created by the disturbance caused by the changing gravitational field, i.e. the changing spacetime geometry. An early example, dating from the late 1960s, was due to the work of Parker (1969), who realized that a general expanding universe would disturb the quantum vacuum and cause particles to be created. (de Sitter space is an exception.) How should we envisage the spontaneous creation of particles from the quantum vacuum as a result of the dynamics of geometry? A helpful heuristic

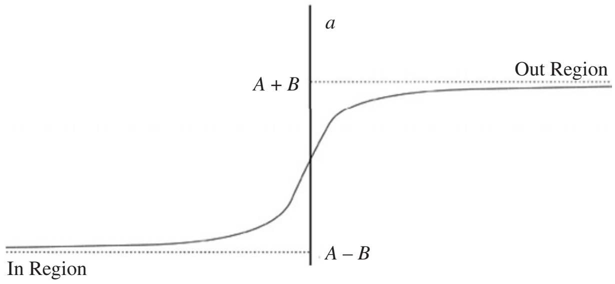

Figure 1. The scale factor of an expanding universe with asymptotically static 'in' and 'out' regions.

model was suggested long ago by DeWitt (1979), who pointed out that the quantum vacuum is in some respects reminiscent of the aether. It can be helpful to think of spacetime as being filled with a type of invisible fluid medium, representing the seething background of vacuum fluctuations. Although the mechanical properties of this medium can be peculiar, and the analogy should not be pushed too far, it is often useful to envisage the 'quantum aether' as possessing a form of viscosity.

To illustrate the concept of vacuum viscosity, consider the effect of an expanding universe. One can think of this either as an external disturbance (the expansion) 'promoting' virtual quanta from the vacuum into real quanta, or as due to the viscosity of the vacuum generating heat as the 'fluid' is expanded. In the case that the expansion is homogeneous and isotropic, this corresponds to bulk viscosity. If the universe expands anisotropically, then shear viscosity of the quantum aether also plays a role, and the particle production is much more prolific (Birrell and Davies 1982, section 5.6). In both cases, the back-reaction of the particle production serves to damp the cosmological motion, and so acts as a genuine viscous drag.

Let me now illustrate mathematically how one may compute the density of created particles in this scenario, using a model first studied by Bernard and Duncan (1977). For simplicity, consider the example of one space dimension, where the expansion is homogeneous and the cosmological scale factor has the form shown in figure 1. Note that in this example there is no big hang. Instead the universe starts out as conventional flat spacetime (Minkowski space), then expands smoothly, and ends up as flat spacetime once more, but with any given initial region of the universe expanded in size by the factor $(A + B)/(A - B)$. The metric for this spacetime is

$$ds^2 = dt^2 - a^2(t)\,dx^2. \tag{6}$$

Here and henceforth I adopt the units $\hbar = c = 1$.

A massless scalar field $\varphi(t, x)$ propagating in this background spacetime satisfies the curved spacetime wave equation

$$\Box\varphi(t, x) = 0 \tag{7}$$

where

$$\Box = (-g)^{-1/2}\partial_\mu[(-g)^{1/2}g^{\mu\nu}\partial_\nu] \tag{8}$$

and g is the determinant of the metric tensor $g_{\mu\nu}$. (I adhere to the sign conventions adopted in Birrell and Davies (1982).) It is possible to solve the equation (7) exactly in a limited number of cases. One example is

$$a(\eta) = [A + B\tanh(\rho\eta)]^{1/2} \tag{9}$$

$$\eta = \int^t dt'/a(t') \tag{10}$$

where A, B, ρ are constants.

It is simpler to work with the so-called conformal time η rather than the cosmic time (proper time) t. A complete set of field modes is then

$$u_k^{\text{in}}(\eta, x) = (4\pi\omega_{\text{in}})^{-1/2}\exp\{ikx - i\omega_+\eta - (i\omega_-/\rho) \times \ln[2\cosh(\rho\eta)]\}\,{}_2F_1(1 + (i\omega_-/\rho), i\omega_-/\rho; 1 - (i\omega_{\text{in}}/\rho); \tfrac{1}{2}(1 + \tanh(\rho\eta))) \tag{11}$$

where ${}_2F_1$ is a hypergeometric function and

$$\omega_{\text{in}} = [k^2 + m^2(A - B)]^{1/2} \tag{12}$$

$$\omega_{\text{out}} = [k^2 + m^2(A + B)]^{1/2} \tag{13}$$

$$\omega_\pm = 1/2(\omega_{\text{in}} \pm \omega_{\text{out}}). \tag{14}$$

Note that in the 'in' region ($t \to -\infty$, $\eta \to t$) these modes reduce to conventional Minkowski space exponential modes:

$$u_k^{\text{in}} \to (4\pi\omega_{\text{in}})^{-1/2}\exp(ikx - i\omega_{\text{in}}\eta). \tag{15}$$

The modes given by equation (11) may be used to define particle states and a Fock space in the Heisenberg picture in the conventional way. In particular, the field φ may be expanded

$$\varphi = \sum_k (a_k u_k^{\text{in}} + a_k^\dagger u_k^{\text{in}*}) \tag{16}$$

and a vacuum state defined by

$$a_k|0_{\text{in}}\rangle = 0. \tag{17}$$

In the 'in' region, $|0_{\text{in}}\rangle$ coincides with the standard definition of a quantum vacuum of normal Minkowski space quantum field theory. However, in the 'out' region, where $t \to +\infty$, the modes given by equation (11) are not simple exponentials, but are more complicated functions of time.

Alternatively, one may find a complete set of modes of the field that reduce to simple exponentials in the 'out' region but not the 'in' region, and use them to define an 'out' vacuum state:

$$u_k^{\text{out}}(\eta, x) = (4\pi\omega_{\text{out}})^{-1/2}\exp\{ikx - i\omega_+\eta - (i\omega_-/\rho) \times \ln[2\cosh(\rho\eta)]\}\,{}_2F_1(1 + (i\omega_-/\rho), i\omega_-/\rho; 1 - (i\omega_{\text{out}}/\rho); \tfrac{1}{2}(1 - \tanh(\rho\eta))) \tag{18}$$

$$u_k^{\text{out}} \to (4\pi\omega_{\text{out}})^{-1/2}\exp(ikx - i\omega_{\text{out}}\eta). \tag{19}$$

These modes are complicated functions of time in the 'in' region. Again, the field may be expanded in terms of these 'out' modes, and an 'out' vacuum state defined:

$$\varphi = \sum_k (b_k u_k^{\text{out}} + b_k^\dagger u_k^{\text{out}*}). \tag{20}$$

$$b_k|0_{\text{out}}\rangle = 0. \tag{21}$$

The significance of the 'out' modes is that they correctly describe the standard definition of vacuum and particle states in the 'out' region (but not in the 'in' region).

The crucial observation is that the 'in' and 'out' modes are *different*. Hence the two vacuum states $|0_{\text{in}}\rangle$ and $|0_{\text{out}}\rangle$ are not the same. That is, the 'in' vacuum contains 'out' particles and vice versa. Since in the Heisenberg picture the state remains unchanged, if we assume the quantum field is in the 'in' vacuum state (i.e. there are no real particles present initially) then there *will* exist particles in the 'out' region. The effect of the period of expansion is to create particles which are detectable in the 'out' region by a standard particle detector. To find out how many one simply solves the wave equation to determine the form of the 'in' modes in the 'out' region, expands them in terms of the 'out' modes, and uses the coefficients to determine the so-called Bogoliubov transformation:

$$u_k^{\mathrm{in}}(\eta, x) = \alpha_k u_k^{\mathrm{out}}(\eta, x) + \beta_k u_{-k}^{\mathrm{out}^*}(\eta, x) \tag{22}$$

where the coefficients are given by

$$\alpha_k = \sqrt{\frac{\omega_{\mathrm{out}}}{\omega_{\mathrm{in}}}} \frac{\Gamma(1 - \mathrm{i}\omega_{\mathrm{in}}/\rho)\Gamma(1 - \mathrm{i}\omega_{\mathrm{out}}/\rho)}{\Gamma(1 - \mathrm{i}\omega_+/\rho)\Gamma(1 - \mathrm{i}\omega_+/\rho)} \tag{23}$$

$$\beta_k = \sqrt{\frac{\omega_{\mathrm{out}}}{\omega_{\mathrm{in}}}} \frac{\Gamma(1 - \mathrm{i}\omega_{\mathrm{in}}/\rho)\Gamma(1 - \mathrm{i}\omega_{\mathrm{out}}/\rho)}{\Gamma(1 - \mathrm{i}\omega_-/\rho)\Gamma(1 + \mathrm{i}\omega_-/\rho)}. \tag{24}$$

Hence the expectation value for the number operator of mode k 'out' particles in the state is

$$\langle 0_{\mathrm{in}}|b_k^\dagger b_k|0_{\mathrm{in}}\rangle = |\beta^2| = \frac{\sinh^2(\pi\omega_-/\rho)}{\sinh(\pi\omega_{\mathrm{in}}/\rho)\sinh(\pi\omega_{\mathrm{out}}/\rho)} \tag{25}$$

which gives the spectrum of created particles for the particular expansion factor equation (9). If this definition of particles seems arbitrary, one may check that if a model particle detector is switched on (slowly) in the 'out' region, it will respond to the 'in' vacuum state in exactly the same way as it would if placed in a conventional quantum state with particle spectrum defined by equation (25). I shall return to the subject of particle detectors in section 5. One of the unsolved problems with this type of calculation is how to define particle and vacuum states when there are no asymptotically static 'in' and 'out' regions. In particular, in the more realistic case where the universe expands from a singular origin, the notion of an initial vacuum is obscure. Over the years there have been many proposals to define 'instantaneous' vacuum states from epoch to epoch as the universe expands (see, for example Grib *et al* 1980), but these definitions have an unappealing ad hoc character. Nor can one use model particle detectors to provide a definition, since these can behave peculiarly when not at rest in Minkowski space, and in any case they suffer from spurious transient effects if switched on abruptly (e.g. at the big bang).

3. Black holes

Some of the most interesting problems of quantum vacuum energy are associated with black holes. Consider, for example, a rotating black hole. According to Einstein's general theory of relativity, such an object should possess a 'magnetic' component that serves to pull an orbiting body around with it, causing it to co-rotate. (This effect is currently being measured for the Earth's rotation in the Gravity Probe B satellite experiment.) The vacuum 'aether' will also be dragged around, but differentially—the effect diminishes with distance. The shearing of the vacuum in this manner creates particles that flow away into the surrounding space, taking angular momentum with them and causing the black hole to slow in its rotation rate. The theory for this rotation radiation was worked out by Starobinski (1973) and Unruh (1974). An unsolved problem is whether a rotating star will also produce such radiation (Matacz *et al* 1993).

Shortly after the Starobinski–Unruh effect was discovered, Stephen Hawking made his famous prediction that non-rotating black holes also emit radiation (Hawking 1975). This time the mechanism is different. An imploding spherical star drags the vacuum aether with it down a black hole, and the resulting heat generated turns out to be precisely thermal. Hawking derived this result by following the sort of procedure I outlined in section 2 for a cosmological model; that is, he decomposed the quantum field in 'in' and 'out' modes, and expanded the 'in' vacuum in terms of 'out' states. Here the 'in' region corresponds to the (almost) flat spacetime prior to collapse, and the 'out' region refers to the spacetime far from the black hole long after the collapse phase is over. Hawking evaluated the Bogoliubov transformation between these two sets of modes, and found for the case of a spherical uncharged black hole of mass M a thermal spectrum with temperature

$$T = 1/8\pi GMk \tag{26}$$

where k is Boltzmann's constant. Hawking concluded that a black hole is not black, but radiates in the same way as a black body. For a solar mass black hole the temperature is a tiny 10^{-8} K, but a curious feature of quantum black holes, clear from equation (26), is that they have negative specific heat. That is, as the hole loses energy, hence mass, it gets hotter. The Hawking effect is therefore unstable, and the black hole radiates faster and faster until its mass approaches zero. The ultimate fate of an evaporating black hole is still a matter of contention.

4. Moving mirrors

A changing gravitational field (i.e. a non-static spacetime) is not the only way to disturb the state of a quantum field. A moving reflecting boundary (mirror) may also create real particles from the quantum vacuum. If a mirror suddenly moves, the information about this change does not reach a distant place until at least the light travel time from the mirror surface to that location, so the vacuum 'aether' in the intervening space is compressed. Vacuum viscosity then leads to heat being generated in the form of particles.

In the case of a one-dimensional mirror (reflecting point) moving in one space dimension, the problem is exactly soluble for a massless scalar field in terms of the energy flow from the mirror (Fulling and Davies 1976, Davies 1977, Davies and Fulling 1977), although the particle spectrum normally requires a numerical treatment. For a mirror trajectory

$$x = z(t) \qquad z = 0, \quad t < 0 \tag{27}$$

the energy flux F is given by

$$F = -(12\pi)^{-1}(1 - v^2)^{1/2}(1 - v)^{-2}\,\mathrm{d}\alpha/\mathrm{d}\tau \tag{28}$$

where v is the mirror velocity, α is the proper acceleration, and τ is the proper time. The actual spectrum of the radiation is computed from a Bogoliubov transformation between 'in' and 'out' modes, as in the cosmological example I considered in section 2. One of the main motivations that Fulling and I had for studying moving mirror radiation in the 1970s was to understand Hawking's black hole radiance effect. In particular, we wanted to know why an imploding star, which leads to an escalating red shift in the radiation emitted from its surface, should produce a steady flux of thermal radiation from the quantum vacuum. We found a direct analogue for the much simpler moving mirror case. The trajectory with asymptotic form

$$z(t) \to -t - A\mathrm{e}^{-2\kappa t} + B \qquad (29)$$

where A, B and κ are constants, provided just such an example. Any radiation reflected from this retreating mirror would reflect back with an exponentially increasing red shift. Nevertheless, at late time the retreating mirror generates a constant flux of energy from its surface, as may be immediately verified using equation (28). Moreover, the Bogoliubov transformation for the moving mirror with trajectory (29) is identical to that for a star that collapses to a black hole (Davies 1977).

In a general application of (28), F need not always be positive. In fact, when the acceleration is increasing to the right, the energy flow to the right is *negative*. This is one example among several scenarios in quantum field theory where negative energy fluxes are possible, and there is a considerable literature examining the implications of this for the second law of thermodynamics (Ford 1978, Davies 1982, Unruh and Wald 1982, Candelas and Sciama 1983, Grove 1988, Ford and Roman 1993). For example, can the entropy of an oven or a black hole be reduced by directing a sustained negative energy beam into it to cool it down? The answer seems to be no. Ford has shown (Ford 1991, Ford and Roman 1993) that the duration of a negative energy flux is normally strictly circumscribed by an uncertainty principle type of inequality which prevents the entropy from going down significantly. However, there are scenarios involving black holes in which the inequality is evaded and there is as yet no general proof that the second law is immune from negative energy effects (Ford *et al* 2002).

Another startling consequence of negative energy fluxes discussed by Ford and Roman (1990, 1992) occurs when the flux is directed at a black hole with maximal electric charge. It is well known that this is a limiting case: if the mass of such a black hole is reduced, the event horizon vanishes, and the black hole is converted into a naked singularity, thus violating cosmic censorship. As a result, the universe is no longer causally secure. Ford and Roman find that the singularity inside a black hole can indeed be briefly exposed by directing a negative energy flux at it, but the situation is rapidly restored by the subsequent inevitable pulse of positive energy. Ford and Roman call this fluctuating horizon 'cosmic flashing'. In effect, the causal influences that might emanate from the singularity are masked by the noise of the fluctuating horizon. But this is not random noise, because the negative energy flux has a predictable form dependent on the mirror motion. It is an open question whether under these circumstances information about the singularity can get out.

Moving mirror radiation is exceedingly feeble unless the accelerations involved are colossal, so there remains doubt over whether it can ever be detected. Attempts have been made to attribute sonoluminescence to moving mirror radiation (Eberlein 1996). This phenomenon occurs when sound is passed through water, causing flashes of light to appear. It is thought that they are generated when small bubbles collapse with enormous rapidity. Treating the bubble as a cavity bounding the electromagnetic vacuum à la Casimir, with the bubble surface playing the role of a partially reflecting mirror, the implosion effectively compresses the quantum vacuum and generates photons. The evidence now seems to be against moving mirror radiation as the principal explanation (Gaitan 1999). In a true Casimir situation, however, involving parallel mirrors, amplification of the moving mirror radiation is possible if one mirror oscillates in resonance with the light travel time across the cavity. Calculations suggest that this form of dynamic Casimir effect may bring moving mirror radiation into the realm of the detectable (Lambrecht *et al* 1996).

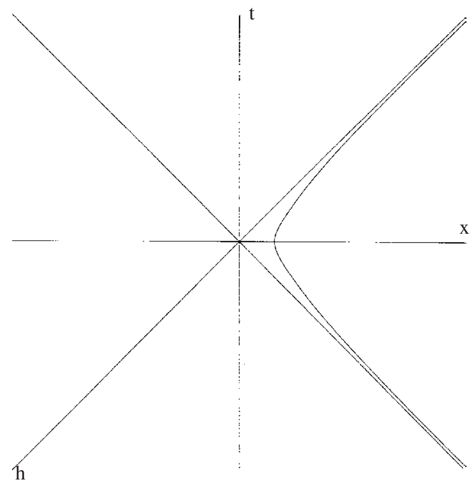

Figure 2. The hyperbola shows the world lines of a uniformly accelerating particle detector.

5. Accelerated systems

One of the most dramatic and much-publicized examples of vacuum noise effects concerns accelerated observers. Some years ago, Unruh and I independently predicted that a uniformly accelerated observer moving through a quantum vacuum would perceive a bath of thermal radiation with a temperature

$$T = \alpha/2\pi k, \qquad (30)$$

where α is the proper acceleration (Davies 1975, Unruh 1976). Unruh, and later DeWitt (1979), showed that a model particle detector in its ground state would respond, when accelerated, in exactly the same way as if immersed at rest in thermal radiation at the temperature given by equation (30). In effect, the quantum vacuum fluctuations, when viewed in an accelerated reference frame, become thermal fluctuations. The situation closely resembles Hawking's black hole radiation effect. In fact, the Bogoliubov transformations are formally identical. The trajectory of a uniformly accelerated observer is a hyperbola in Minkowski space, shown in figure 2. The asymptotes lie along the light cone through the origin, which play a role analogous to the event horizon (h) in the black hole.

An interesting question concerns the conservation of energy of an accelerated detector. Since the detector becomes excited, it must gain energy. On the other hand, from the standpoint of a stationary observer, the quantum field is initially in a vacuum state. Thus, any transition of the detector to an excited state must involve a transition in the field too, and in first-order perturbation theory the only possible transition from the vacuum state is the emission of a quantum. So in the frame of the detector energy is absorbed, but in the un-accelerated frame it is emitted!

The paradox is resolved when it is realized that the initial state of accelerated detector + vacuum is not a total energy eigenstate. In an individual transition, therefore, energy is not conserved. Consider an ensemble of accelerated two-level detectors that have their energies measured at the end of some time interval. Since in the accelerated frame there is an apparent thermal bath, some detectors will be excited, and others will be in the ground state, with energies distributed according to the usual Boltzmann factor with temperature determined by equation (30). For those detectors that are excited, the expectation value of the energy will have gone up, but for the unexcited ones the expectation value will have gone down. Combining these possibilities, it turns out that the total energy expectation value remains unchanged (Grove 1986).

Linear acceleration is the most striking example, but almost any sort of acceleration will produce a similar effect, i.e. the perception of a bath of radiation in what for an inertial observer is a vacuum state. Generally, however, the spectrum of the radiation is non-thermal. One case of interest is uniform rotation (Letaw and Pfautsch 1980). The energetics here are different from the case of linear acceleration. In the non-rotating frame, when a rotating detector becomes excited it emits a quantum into the field that carries angular momentum away and serves to damp the motion of the detector. This is therefore another vacuum friction effect. This time, however, there is a net energy loss from the detector to the field. To sustain the rate of angular motion, energy must be continuously fed into the system (Davies *et al* 1996).

The magnitude of 'acceleration radiation' is disappointingly feeble. An acceleration of 2.5×10^{22} cm s^{-1} is needed to generate an effective temperature of just 1 K. Nevertheless Bell and Leinaas (1983, 1987) claim to have seen a positive effect in the spin depolarisation of electrons in a particle accelerator ring.

6. Decoherence and dissipation

My final example of vacuum friction is much closer to experimental study than the foregoing. Moreover, it involves frictional forces in a very familiar context. Imagine a Casimir situation in which one plate moves, not orthogonal to the plane of the plates as I discussed at the end of section 4, but in the plane of the plate, i.e. one plate 'slides across' the other. If the plates were perfectly conducting, nothing would happen because of Lorentz invariance; put simply, the plates would not know they were in relative motion because they lack any markers. However, real plates are not perfect conductors/reflectors. They are composed of dissipative materials. As a result, there will be a frictional force experienced that acts to damp the relative motion. Its magnitude has been calculated by Pendry (1997), and an experimental test carried out by Chen *et al* (2002). Since the region between the plates is a quantum vacuum, the friction is entirely a vacuum effect. No real photons are involved.

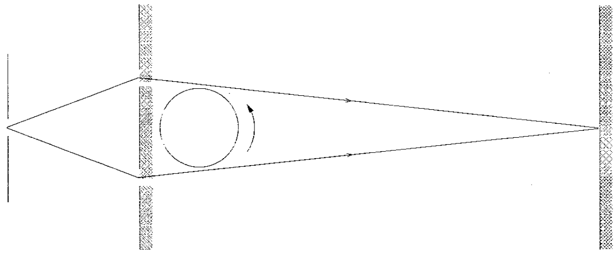

Figure 3. Modified atom interference experiment. The rotating cylinder positioned near the slits exerts a differential vacuum viscosity on the passing atoms, which serves to both shift the interference pattern and decohere it.

This sort of phenomenon is at its most striking in the case of a single atom moving parallel to, but some distance from, an imperfectly conducting plate. The atom also experiences a velocity-dependent damping force due to vacuum friction. The kinetic energy of the atom appears as heat in the plate; virtual photons transfer the energy from the atom to the plate. One way to envisage the phenomenon is as follows. An atom located near a reflecting surface sees an image of itself and will experience an effective van der Waals attractive force. In the case of a transversally moving atom, the image moves parallel to it, but the dissipation in the plate causes the image to lag slightly behind the atom. As a result, the effective attractive force between the atom and the image has a small component parallel to the plate that acts to retard the atom's motion. This example of 'friction-at-a-distance' is reminiscent of the Moon's motion around the Earth. Friction heats the tidal bulge in the rotating earth, and the resulting lag in the bulge relative to the moon's motion creates a retarding force that causes the moon to lose energy and slow down in its orbit.

It is possible that vacuum friction in this form may be measurable experimentally. For example, if an atom is dropped vertically down the centre of a metal cylinder, it should reach a terminal velocity due to vacuum friction with the material of the cylinder. In practice, the atom's vertical trajectory would be unstable due to its attraction to the surface of the tube, so an alternative strategy is to exploit a discovery made by Ford and Svaiter (2002) in which a parabolic mirror is used to focus vacuum fluctuations. Ford and Svaiter find the existence of a small restoring force that has the effect of drawing an atom towards the axis of the parabola. If the mirror was placed vertically, the focused vacuum would have the effect of creating a 'virtual pipe' or waveguide down which the atom could fall. Another reflecting surface (e.g. a cylindrical pipe with curvature away from the axis) brought into close proximity of the axis would then provide the source of vacuum friction. The fall of the atoms could be monitored with lasers.

Vacuum friction may also be observable in interference experiments. Suppose a standard two-slit experiment is performed with atoms, but a rotating metal cylinder is inserted between the slits (see figure 3). There should then be a shift in the interference pattern, because the atoms moving through one slit will be accelerated while those moving through the other slit will be retarded. Furthermore, since the cylinder material will be dissipative, there will be some loss of phase coherence, which will serve to reduce the overall degree of interference. By adjusting the conductivity and rotation rate of the cylinder, these two effects can be independently tuned. There is considerable interest in the study of decoherence in quantum mechanics (Zurek 1991). A dissipative environment provides a very strong source of decoherence, but the relationship between the decoherence time and the dissipation time is a subtle one. The experiment proposed above could be used to investigate the interweaving of these two fundamental effects, one controlling the emergence of classicality in the universe, the other the emergence of an arrow of time.

7. Conclusion and future experiments

The Casimir effect and other vacuum energy phenomena assume a much wider significance when time-dependence is taken into account. I have argued that a heuristic way to regard dynamic vacuum energy effects is by appealing to a sort of vacuum friction. This leads to mechanical back-reaction effects, such as the slowing of rotating black holes or the viscous drag between moving plates. It also leads to particle creation from the vacuum.

A largely unexplored area of research for the dynamic Casimir effect is in the realm of biophysics. The living cell is replete with nanomachines, including pumps, rotors, linear motors, pincers and nanotubes. Casimir-type forces will be very considerable for these systems. A study of the quantum properties of cellular components is likely to prove very fruitful (Davies 2004). The cell membrane is a uniform slab of dielectric material that will produce changes in the quantum vacuum energy both within the membrane and in the vicinity of its surfaces. These changes will affect the energy levels of molecules close to the surfaces, which will in turn affect the physical and chemical properties of these substances. It is known that the properties of water change markedly close to cell membranes in a manner that may be crucial for biological function. Part of this alteration may be attributable to quantum vacuum energy or vacuum friction.

Despite their peripheral theoretical origins, quantum vacuum effects have turned out to be of immense importance, and probably play a key role in shaping the universe. In particular, they seem to have provided the all-important primordial density fluctuations that triggered the growth of galaxies, a step that led eventually to life. But life consists of cells that are collections of nanomachines in the vicinity of dielectric surfaces. It is an arresting speculation, therefore, that quantum vacuum effects may not only hold the key to the universe; they may also hold the key to life itself.

Acknowledgments

I am grateful to Larry Ford and Claudia Eberlein for helpful discussions.

References

Bell J and Leinaas J 1983 *Nucl. Phys.* B **212** 131
Bell J and Leinaas J 1987 *Nucl. Phys.* B **284** 488
Bernard C and Duncan A 1977 *Ann. Phys.* **107** 201
Birrell N D and Davies P C W 1982 *Quantum Fields in Curved Space* (Cambridge: Cambridge University Press)
Bunch T S and Davies P C W 1978 *Proc. R. Soc.* A **354** 59
Candelas P and Sciama D 1983 *Phys. Rev.* D **27** 1715
Chen F, Mohideen U, Klimchitskaya G L and Mostepanenko V M 2002 *Phys. Rev. Lett.* **88** 101801
Davies P C W 1975 *J. Phys. A: Math. Gen.* **8** 609
Davies P C W 1977 *Proc. R. Soc.* A **356** 237
Davies P C W 1982 *Phys. Lett.* B **113** 393
Davies P C W 2004 *BioSystems* **78** 69
Davies P C W, Dray T and Manogue C 1996 *Phys. Rev.* D **53** 4382
Davies P C W and Fulling S A 1977 *Proc. R. Soc.* A **354** 529
DeWitt B S 1979 Quantum gravity: the new synthesis *General Relativity: An Einstein Centenary Survey* ed S Hawking and W Israel (Cambridge: Cambridge University Press)
Eberlein C 1996 *Phys. Rev. Lett.* **76** 3842
Ford L H 1975 *Phys. Rev.* D **12** 2963
Ford L H 1978 *Proc. R. Soc.* A **364** 227
Ford L H 1991 *Phys. Rev.* D **43** 3972
Ford L H, Helfer A D and Roman T A 2002 *Phys. Rev.* D **66** 124012
Ford L H and Roman T A 1990 *Phys. Rev.* D **41** 3662
Ford L H and Roman T A 1992 *Phys. Rev.* D **46** 1325
Ford L H and Roman T A 1993 *Phys. Rev.* D **48** 776
Ford L H and Svaiter N F 2002 *Phys. Rev.* A **66** 062106
Fulling S A and Davies P C W 1976 *Proc. R. Soc.* A **348** 393
Gaitan F 1999 *Phys. World* **12** (3) 20
Gong Y and Duan C-K 2004 *Mon. Not. R. Astron. Soc.* **352** 847
Grib A, Mamaev S and Mostepanenko V 1980 *Quantum Effects in Strong External Fields* (Moscow: Atomizdat)
Grove P 1986 *Class. Quantum Grav.* **3** 801
Grove P 1988 *Class. Quantum Grav.* **5** 1381
Guth A 1981 *Phys. Rev.* D **23** 347
Hawking S W 1975 *Commun. Math. Phys.* **43** 199
Lambrecht A, Jaekel M and Reynaud S 1996 *Phys. Rev. Lett.* **77** 615
Langlois D 2004 Inflation, quantum fluctuations and cosmological perturbations *Preprint* hep-th/0405053
Letaw J and Pfautsch J 1980 *Phys. Rev.* D **22** 1345
Linde A 1990 *Inflation and Quantum Cosmology* (Boston, MA: Academic)
Matacz A, Ottewill A and Davies P C W 1993 *Phys. Rev.* D **47** 1557
Parker L 1969 *Phys. Rev.* **183** 1057
Pendry J B 1997 *J. Phys.: Condens. Matter* **9** 10301
Sakharov A D 1967 *Dokl. Akad. Nauk Ser. Fiz.* **177** 70
Sakharov A D 1968 *Sov. Phys.—Dokl.* **12** 1040 (Engl. transl.)
Schwarz D J and Terrero-Escalante C A 2004 *J. Cosmol. Astropart. Phys.* JCAP08(2004)003
Starobinski A 1973 *Sov. Phys.—JETP* **37** 28
Unruh W G 1974 *Phys. Rev.* D **10** 3194
Unruh W G 1976 *Phys. Rev.* D **14** 870
Unruh W G and Wald R 1982 *Phys. Rev.* D **25** 942
Zurek W 1991 *Phys. Today* **44** 33