Characterisation of a three-dimensional Brownian motor in optical lattices

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Abstract. We present here a detailed study of the behaviour of a three dimensional Brownian motor based on cold atoms in a double optical lattice [P. Sjölund \textit{et al.}, Phys. Rev. Lett. \textbf{96}, 190602 (2006)]. This includes both experiments and numerical simulations of a Brownian particle. The potentials used are spatially and temporally symmetric, but combined spatiotemporal symmetry is broken by phase shifts and asymmetric transfer rates between potentials. The diffusion of atoms in the optical lattices is rectified and controlled both in direction and speed along three dimensions. We explore a large range of experimental parameters, where irradiances and detunings of the optical lattice lights are varied within the dissipative regime. Induced drift velocities in the order of one atomic recoil velocity have been achieved.

PACS. 32.80.Lg Mechanical effects of light on atoms, molecules, and ions – 05.40.Jc Brownian motion – 32.80.Pj Optical cooling of atoms; trapping

1 Introduction

Brownian motors (BMs) are small scale engines that convert random fluctuations into deterministic work. In order to realise such a device, symmetry has somehow to be broken (Curie’s principle \cite{1}), and the system has to be out of thermodynamical equilibrium \cite{2}. Much effort has been invested in studying the underlying mechanisms of BMs operating with the ratchet effect \cite{2,3,4,5}. The ratchet mechanism consists in breaking the spatial and/or the temporal inversion symmetry of the system so that directed transport emerges, using fluctuations as the relevant input. The paradigmatic device is Feynman’s famous ratchet and pawl machine \cite{6}, based on an idea of von Smoluchowski \cite{7}.

During the past years there has been an extensive body of work on BMs and ratchets, several realised in systems of periodic arrays of ultracold atoms, known as optical lattices (OLs) \cite{8,9,10,11,12}. Optical lattices are periodic arrays of micro-traps created by the interference between two or more laser beams \cite{13,14}. It has been shown that ultracold atoms stored in OLs can be controlled and manipulated with a very high degree of precision and flexibility and they are routinely used for example in studies of Bose-Einstein condensation and applicable for quantum state manipulation \cite{15,16}.

In this paper, we present an experimental and numerical study of fluctuations rectified into directed motion of Brownian particles in three dimensional OLs \cite{9}. This works by purely optical fields, where ultracold atoms are optically pumped between two state-dependent and spatially overlapped OLs \cite{17,18}, both spatially and temporally symmetric. These OLs are coupled by optical pumping, with strongly asymmetric transfer rates, via the vacuum field reservoir, which, together with a shifted relative spatial phase between the OLs, causes atoms to be propelled in a controllable direction.

A general understanding of the dynamics of BMs is of fundamental interest. In biology, directed transport and molecular motors are both driven by stochastic motion, thus providing the energy input for the ability of living cells to generate motion and forces, \textit{e.g.}, for mobility, contraction of muscles or material transport, and are in themselves a fundamental mechanism of the origin of living cells. For example, biological motor proteins which move along linear filaments can be described by stochastic models coupled to chemical reactions \cite{19}. So-called ratchet models further explain the generation of directed motion on the microscopic level out of Brownian motion. The general concept of our BM, based on the idea of \cite{20}, may be transferable into fields such as chemistry and biology as a tool for studies and control of, \textit{e.g.}, molecular motors, investigations of intra-cell motion and possible studies of Brownian motion in biological membranes \cite{21}.

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By letting the relative spatial phase between the two potentials be slightly shifted, combined with the pronounced difference in the transfer rates, the spatiotemporal symmetry of the system is broken and a directed motion is obtained. The concept can therefore be adapted to induce directed motion in any direction in three dimensions since the same physical process will function in any direction in a three-dimensional case. Because we use symmetric potentials, the direction of the induced drift can be reversed by a proper choice of the relative spatial phase shift. This is to be contrasted with the usual use of a ratchet potential (e.g., using a sawtooth potential), where the direction is fixed by the shape of the potential.

3 Experiment

3.1 Experimental setup

The BM is realised, using cold caesium atoms in a double optical lattice (DOL) [17,19,23]. A magneto-optical trap (MOT) is loaded from a chirped-slowed atomic beam, produced in a thermal source which fills the MOT with approximately $10^8$ atoms at a peak density of $10^{11}$ cm$^{-3}$. Approximately 80% of the initial number of atoms from an optical molasses of a few microkelvin is loaded into the DOL with a filling fraction of about 0.05 atom per site. After loading, the OL light is left on for a chosen time $\tau$. To avoid spurious spatial drifts of the atomic cloud caused by light pressure, great care has been taken to balance the laser beam power [28]. Also, it is crucial that the relative spatial phase is kept constant.

This three-dimensional DOL is composed of two spatially overlapped DOLs with identical topography, but which can be controlled individually in terms of both well depths and relative spatial phase. Two different beams of frequencies $\omega_A$ and $\omega_B$ are split into four branches with equal power. This lattice geometry is a generalisation of the 1D lin-in-lin configuration to 3D [14], see figure 2a. We use frequencies that are near-resonant with the $1S_{1/2}$ line of caesium at 852 nm reaching the excited fine structure level $6p\,^2P_{3/2}$. Each optical lattice operates from a different hyperfine ground state of the $6s\,^2S_{1/2}$ level. One operates within the closed $F_g = 4 \rightarrow F_e = 5$ transition, trapping atoms in the $F_e = 4$ ground state. The second lattice traps atoms in the $F_g = 3$ ground state, operating in the open $F_g = 3 \rightarrow F_e = 4$ transition, see figure 2b. These resonances are much narrower than the difference in the laser frequencies, which allows us to address the two lattices independently. At the same time, the wavelengths are close enough, and the trapped atomic cloud is small enough (typically 1–2 mm in diameter), to ensure that the periodicity is the same within the sample volume (e.g., it takes about 3.3 cm for the lattices to phase out by $\pi$ in the horizontal $x$ and $y$-directions). Moreover, we can individually control the optical pumping rates between the lattices to some extent [13] by tuning the irradiances $I_A$, $I_B$ and the detunings $\Delta_A$, $\Delta_B$ of the lattice beams, which in turn changes the well depths, optical pumping rates, diffusion...
Fig. 2. a) Double optical lattice beam geometry, realised by four branches, where each branch contains two spatially overlapped laser beams A and B. $e_1$ to $e_4$ are the polarisation vectors perpendicular to the plane of propagation. b) Schematic energy level diagram: Two ground states $F_e = 4$ and $F_e = 3$ are connected to the excited $F_e = 5$ and $F_e = 4$ states respectively by two laser fields B and A detuned by $\Delta_A$ and $\Delta_B$. $/2\pi = 5.22$ MHz is the natural linewidth of the transitions.

and friction, and consequently, also the behaviour of our BM.

By letting the two lattices be out of phase, e.g., in the vertical $z$-direction, while simultaneously having an asymmetric transfer rate between the lattices, the motion of atoms is channelled in that particular direction [9]. This process can be controlled in all three dimensions, with regard to both speed and direction, by a proper choice of the relative spatial phases $\varphi_x$, $\varphi_y$, $\varphi_z$, irradiances $I_A$ and $I_B$ and the detunings $\Delta_A$ and $\Delta_B$ of the laser light.

The control of the relative spatial phase shifts between the lattices in three-dimensions is achieved by extending or shortening the lattice beam branches 1, 2, 3 and 4, see figure 2a. A relative phase shift only in the horizontal $x$-direction corresponds to an extension of branch 2 and a shortening of branch 4 by the same amount, or vice versa. In $y$, an extension of branch 1 and shortening of branch 2 by the same amount, or the other way around, will cause a shift. For a phase shift only in the vertical $z$-direction, a simultaneous extension or shortening of both branches 1 and 2 or branches 3 and 4 by the same amount will work. Combining these, an arbitrary combination of $\varphi_x$, $\varphi_y$ and $\varphi_z$ can be set. The relative spatial phases between the lattices are experimentally determined by the kinetic temperature, measured by a ballistic time-of-flight (TOF) technique [29], when changing the relative spatial phase, see figure 3. A Gaussian fit to the TOF-signal also provides the peak arrival time of the distribution and the number of atoms by the area under the fitted curve.

3.2 Calibration of the relative spatial phase by temperature measurement

The temperature dependence is shown in figure 3 for a relative spatial phase shift in the vertical $z$-direction. A 3.3 cm extension or shortening of the beam branches correspond to a relative spatial phase shift of $2 \pi$. This is repeated for the other directions in order to find the origin of all the relative spatial phases. An simplified explanation to the temperature variations [17] is an increase in the number of scattered photons when a $\sigma^+$ point in lattice A overlaps with a $\sigma^-$ point in lattice B. Figure 4b, shows a semi-classical model where a $\sigma^+$ point in lattice A overlaps with a $\sigma^+$ point in lattice B. If an atom is optically pumped from lattice B to lattice A, it will end up in the lowest light shift potential (due to conservation of angular momentum), and then return to lattice B. In this process, an atom only scatters a few photons. The opposite situation, where a $\sigma^+$ point overlaps with a $\sigma^-$ point is shown in figure 4a. Here, an atom in lattice B that is pumped to lattice A will end up in the least light shifted anti-trapping potential. The atom will slide down the potential, gaining kinetic energy and, as it is close to a $\sigma^-$ point it will scatter photons, which pumps the atom to the lowest potential. It is then pumped back to lattice B in a similar process. Thus, in this case the number of scattered photons is considerably higher, which result in heating [17,18]. In a more realistic picture, there is a number of Zeeman sub-levels, all corresponding to different light shift potentials. For a $\sigma^+/\sigma^-$ site overlap, photon scattering increases since atoms are optically pumped between a range of different $M$-states.

3.3 Measurement of drift velocity by TOF

In figure 5 three TOF signals are shown, all for $\tau = 350$ ms. The middle one, indicated by (0), is when the two OLS

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1 This is done by manually adjusting the distance between two facing prisms, mounted on linear translation stages in each of the four branches.
Fig. 4. a) If two σ⁺ points in lattice A (ground state |gA⟩) and lattice B (ground state |gB⟩) overlap, atoms can be transferred between the lattices with minimised light scattering. b) In the opposite case, when σ⁻ overlaps with σ⁺, heating effects are enhanced.

Fig. 5. (Colour online) Normalised TOF-signals with arrival time t as the horizontal axis, for relative spatial phases that generates maximum induced drifts in the upward (+z) and downward (−z)-direction. The middle one (0) shows a TOF signal for no induced drifts when the optical lattices are in phase. All three TOF signals are normalised to make comparison easier and the signals are averages of four drops.

are in phase and no induced drift is present. Those indicated by (+z) and (−z) are when φz is set to generate a maximum drift, either upwards or downwards. A more direct technique, and a complement to the TOF technique, is imaging the shadow of the atoms, transiently illuminated by a weak resonant probe beam, on a CCD detector, see section 3.4. The images provides the centre-of-mass position of the atomic cloud in the xz plane, determined by Gaussian fits to the images after an arbitrary lattice time τ.

In order to experimentally investigate the behaviour of the BM, we change the parameters ΔA and ΔB, covering the frequency span between the excited $F_e = 3$ to the $F_e = 4$ states, and the $F_e = 4$ and the $F_e = 5$ states respectively. $I_A$ and $I_B$ are set independently up to about 20 mW/cm². To characterise the drift velocity in the vertical z-direction ($v_{dz}$), we tune the parameters ($Δ_A$, $Δ_B$, $I_A$, $I_B$) in a number of steps, all for a fixed lattice time τ of 350 ms. For each set of parameters, the relative spatial phase in the vertical z-direction ($φ_z$) was incrementally changed.

Fig. 6. Induced drift velocities $v_{dz}$ as a function of relative spatial phase $φ_z$ for three different sets of parameters. Filled circles: $Δ_B = 40Γ$, filled squares: $Δ_B = 25Γ$ and circles: $Δ_B = 10Γ$, for a lattice time τ of 350 ms, $Δ_B = 14Γ$, $I_A = 0.8$ mW/cm² and $I_B = 1.2$ mW/cm² for all cases.

In about 50 steps, covering slightly more than 2π, while $φ_x$ and $φ_y$ are kept at zero, meaning that only vertically induced drift was present, either upwards or downwards. In total, we measured about 14000 velocity distributions, due to the large parameter space. For each change of $φ_z$, five TOF-signals were averaged to increase the signal-to-noise ratio and Gaussian fits were performed to read out the peak arrival time of the TOF signals.

With a constant drift velocity $v_0$, the vertical drift velocity $v_{dz}$ is determined from the peak arrival time $t$ of the atoms at the TOF probe by

$$v_{dz} = \frac{gt^2 - 2l}{2(t + τ)},$$

where $l$ is the vertical distance that the atoms fall down to the probe and $g$ is the gravitational acceleration. In figure 6, three different curves are plotted, showing $v_{dz}$ as a function of $φ_z$ for three different values of $Δ_B$, while $Δ_A$, $I_A$ and $I_B$ are kept fixed. Here, a clear variation of $v_{dz}$ is evidenced for different values of $Δ_B$. In figure 7, $Δ_B$ is plotted as a function of $v_{dz}$, for a relative spatial phase that generates a maximum $v_{dz}$ in the upward (+z)-direction, with $Δ_A$, $I_A$ and $I_B$ fixed for a lattice time τ of 450 ms. The maximum drift velocity clearly increases when $Δ_B$ increases and approaches the $F_e = 4$ state. Continuing the increase of $Δ_B$ close to the $F_e = 4$ state results in an abrupt decrease of the TOF-signal due to resonant trap losses. However, changing $Δ_A$, $I_A$ and $I_B$, while keeping $Δ_B$ fixed, affect $φ_z$ less significantly and in a more ambiguous way compared with changing $Δ_B$, as can be seen in figure 8. With the large parameter space and with the high degree of coupling between the different parameters, it is difficult to extract any unambiguous trends from the data. The behaviour of this type of BM is complex and since it is working in the dissipative regime, a strong coupling between friction, heating and dissipation is present. Changing any parameter, simultaneously changes the fric-
Fig. 7. Maximum vertical drift velocity $v_{d,z}$ as a function of $\Delta B$, where $\Delta A = 33F$, $I_A = 4.17$ mW/cm$^2$, $I_B = 2.71$ mW/cm$^2$ and the lattice time $\tau = 450$ ms. The $F_k = 4 \rightarrow F_a = 4$ resonance occurs at $\Delta B = 48F$.

Fig. 8. Maximum drift velocity $v_{d,z}$ in the (+z)-direction for different sets of parameters, plotted as a function of $I_A$. Meanwhile, $\Delta_B = 40F$ is kept fixed during a lattice time $\tau = 350$ ms.

3.4 Evidence for horizontal and vertical drifts

Our lattice structures are periodic in three dimensions. We can adjust the relative spatial phases $\varphi_z$, $\varphi_y$ and $\varphi_x$ at will. Thus, the Brownian motor works also horizontally, and indeed in an arbitrary direction. To confirm this we measure the position of the atomic cloud in the $xz$-plane as a function of the lattice time $\tau$. This is done by imaging the shadow of the atomic cloud, transiently illuminated by a weak probe beam, on a CCD-detector. Figure 9 shows false colour images of the atomic cloud for different successive lattice times. An induced drift is evident both along $z$ and $x$, and as well as in the diagonal $xz$-direction for an appropriate choice of $\varphi_z$ and $\varphi_y$ ($\varphi_x = 0$) which generates the largest drift velocities while keeping $\Delta A$, $\Delta B$, $I_B$ and $I_A$ fixed. To confirm the induced drift dependence along $x$ as a function of $\varphi_x$, we determine for each change in $\varphi_x$ the centre-of-mass $x$-position by a Gaussian fit to the images, while $\varphi_z$ and $\varphi_y$ where kept fixed at zero. For comparison reason with TOF results in the $z$-direction, the same lattice time $\tau = 350$ ms was used for one typical set of parameters, see figure 10.

4 Numerical simulations

In order to understand the qualitative behaviour of our BM, we have performed simulations, using a simple model displaying the basic characteristics of our system. We consider a classical Brownian particle which can be in one of two internal states, indexed by $j$, and interacting with state-dependent external potentials $U_j$. In analogy to the experimental setup, section 3, we chose $U_B$ to be the light shift potential resulting from laser field B (see figure 2), corresponding to the lowest adiabatic optical potential $\Gamma$ for the $F_k = 4 \rightarrow F_a = 5$ transition of the 4-beam linear configuration. Consequently, $U_A$ is chosen as the
Fig. 10. Experimental results showing the drift velocity \( v_{dA} \) as a function of relative spatial phase along \( x \). Here, \( \Delta_B = 40 \Gamma \), \( \Delta_A = 36 \Gamma \), \( I_B = 6.13 \text{ mW/cm}^2 \) and \( I_A = 6.67 \text{ mW/cm}^2 \) for a lattice time \( \tau \) of 350 ms. The error bars come from the uncertainty in the Gaussian fits used to determine the centre-of-mass \( x \)-position to the images.

Fig. 11. (Colour online). One dimensional cuts \( U(0,0,z) \) (blue) and \( U(x,0,0) \) (red) of the potential surfaces used in the simulation, expressed as a function of the light shift \( |\Delta_0| \) \cite{14}, where \( k \) is the absolute value of the wavevector of the laser field. Solid line: lowest adiabatic potential of the \( F_6 \) = 4 \( \rightarrow \) \( F_0 \) = 5 transition; dashed line: lowest adiabatic potential of the \( F_6 \) = 3 \( \rightarrow \) \( F_0 \) = 4 transition.

Fig. 12. Results from numerical simulations of the Brownian motor mechanism for the vertical \( z \) (full line) and horizontal \( x \) (dashed line) directions. The drift velocity \( v_d \) is plotted as a function of the relative spatial phase \( \varphi = \varphi_A - \varphi_B \), varied independently along \( z \) or \( x \).

case \( j \) as A or B, reads \cite{20, 30}

\[
[\partial_t + v \partial_x] W_j - \partial_v [v + \nabla U_j(x) + D_{v,j}(x) \partial_v] W_j = \gamma_{j' \rightarrow j}(x) W_{j'} - \gamma_{j \rightarrow j'}(x) W_j
\]

for \( j' \neq j \), with time expressed in units of the inverse of the friction coefficient \( \alpha \) and position in units of the inverse angular wave vector \( k \) of the laser light, such that all variables are dimensionless. In addition, we scale the potentials by a factor of \((45/88)\mathcal{A}\) such that the adiabatic potential along \( z \), \( U(0,0,z) \) shown in figure 11, has a depth of \( \approx 2\mathcal{A} \), allowing a direct comparison with the sine potential amplitude \( \mathcal{A} \) used in ref. \cite{20}. We assume that the mean kinetic energy of the system is smaller than the depth of the trapping potentials and that the typical frequency of the oscillations in the potential wells is larger than the rate of energy damping, so that the dynamics falls into the low-damping regime. Motion is restricted to 2D (\( y = 0 \)) and we observe the average velocity \( v_d \) in both \( x \) and \( z \) directions as a function of the relative phase shift \( \varphi \) in either of these directions. Unless noted otherwise, we use arbitrary position-independent values of \( D_v \) and \( \gamma \), with \( \gamma_{A \rightarrow B} = 3\gamma_{B \rightarrow A} = 7.5 \) (i.e., an atom spends 75% of its time in the long-lived lattice B).

Figure 12 shows the dependence of the Brownian motor mechanism on the relative phase \( \varphi \) between the two potentials, with \( \mathcal{A} = 200 \) and \( D_v = 75 \) in both states. The overall shape of the curve is similar to what was obtained for a sine potential in ref. \cite{20}, with a greater drift velocity for a phase shift along \( x \), due to the deeper potential along that direction. Indeed, figure 13 shows that the drift velocity is proportional to the potential depth, up to a certain value where the BM effect saturates. Similarly, changing the diffusion constant \( D_v \) affects the BM differently depending on the internal state of the atom, as seen in figure 14. Increasing the diffusion in both lattices, or

\( A \) coefficient of 44/45 appears in the formula for the potential from the Clebsch-Gordan coefficients for a \( F_6 = 4 \rightarrow F_0 = 5 \) transition.
only in the long-lived one, results in a decrease of the drift velocity, due to the added noise. Conversely, a greater diffusion in the short-lived lattice slightly increases the drift velocity, because the atom can take advantage of the increased Brownian motion without too much adverse effect from isotope diffusion.

Fig. 13. Dependence of the magnitude of drift velocity on the depth of the potentials. Circles: identical depth $A_B = A_A$; squares: varying depth of the long-lived lattice $A_B$ for $A_A = 200$; triangles: varying depth of the short-lived lattice $A_A$ for $A_B = 200$. $\varphi_+ = 2\pi/3$ and $D_v = 75$ in all cases.

Fig. 14. Dependence of the magnitude of drift velocity on the diffusion constant. Circles: identical diffusion $D_{v,A} = D_{v,B}$; squares: varying diffusion in the long-lived lattice $D_{v,B}$ for $D_{v,A} = 75$; triangles: varying diffusion in the short-lived lattice $D_{v,A}$ for $D_{v,B} = 200$. $\varphi_+ = 2\pi/3$ and $A = 200$ in all cases.

5 Discussion

There are two ambient forces that may affect the experimental results. One is due to earth magnetic field, from which a slight Zeeman shift could be introduced. This is cancelled by a B-field compensation in the experimental setup. The second force is due to gravity. It will introduce a slight tilt of the lattice potentials, which may result in an increased tunnelling probability in one direction. This effect is small, and for potential depths used in the experiment of 100–300$E_r$. The potential energy $U_B = mgh$ contribution due to gravity becomes $\sim 10^{-3}E_r$, where $m$ is the atomic mass, $g$ is the gravitational constant and $h = \lambda/\sqrt{2}$ is the distance between two $\sigma^+\sigma^-$-sites in the vertical $z$-direction ($\lambda = 852.3$ nm). This clearly indicates that this contribution is negligible compared to the potential depth.

While two periodic potentials are used in the model, in reality the atoms have magnetic substates, each leading to a different light shift potential. A further difficulty with our BM is that, since it works in a dissipative regime, the parameters, potential depths, magnitude of the diffusion, and the transition rates between the two lattices all depend on the irradiances and the detunings of the OL lasers. Changing any parameter results in a different friction, diffusion and as well as dissipation. Therefore, one of the main features of this paper is to investigate how this type of Brownian noise rectifier varies under different conditions. The complexity of the system makes it difficult to directly relate the experimental findings to the simple classical model used in the simulation. Quantum mechanical simulations in terms of parameters directly controllable in experiments, using the full level structure, are under development.

Despite these difficulties, carefully investigations of the parameter space, $\Delta_A$, $\Delta_B$, $I_A$ and $I_B$, allow some general conclusions to be drawn. The most striking feature is the drastic increase of the effectiveness of the BM when $\Delta_B$ is increased. As $\Delta_B$ is increased, the frequency of the light will get closer to the $F_g = 4 \rightarrow F_e = 4$ transition, and hence the pumping rate $\gamma_{B \rightarrow A}$ increases. Classical simulations [20] confirm that there is an optimal ratio between pumping rates, as long as the inequality remains large. Changing $\Delta_B$ will also change the pumping rate between the sublevels within the lattice B manifold. Therefore, the potential depth, friction and diffusion within lattice B will also be modified, both in magnitude and position dependence, in a non-straightforward way [22].

However, in the limit $\gamma_{B \rightarrow A}/\gamma_{A \rightarrow B} \rightarrow 0$ or $\infty$ the system is effectively reduced to a single optical lattice, and we would expect any BM effect to vanish. Therefore it is clear that the optimal BM effect must be achieved for some finite ratio $\gamma_{B \rightarrow A}/\gamma_{A \rightarrow B}$, as has indeed been confirmed by classical simulations [20]. This is clearly consistent with the data in Fig. [4]. Due to experimental difficulties when laser B is tuned too close to level A we have not been able to extend our data to larger $\Delta_B$ in order to investigate if the drift velocity falls off again. Classical simulations indicate that this will happen for $\gamma_{B \rightarrow A}/\gamma_{A \rightarrow B}$ larger than
0.1-0.2. Although we do not know the exact relation between $\gamma_{A\rightarrow B}/\gamma_{A\rightarrow B}$ and $\Delta_B$ (which will also depend on other parameters) measurements of the relative populations of the two lattices indicate that this value has not been reached[18].

For quantitative analyses, the model using only two potentials, while ignoring the internal level structure of the two lattices, is not sufficient. This is evident since features like the ones at $\varphi = 0$ and at $\varphi = 2\pi$ for the directional variation of the shape of the velocity curve in figure 12 do not show up in the experimental graphs (see figure 6). It is also expected that the position dependence of pumping, friction and diffusion coefficients will greatly affect the shape of the curve.

6 Conclusion

In summary, we have demonstrated a Brownian motor working in three dimensions with a controllable speed. Induced drift velocities in the order of one recoil velocity have been achieved. We have showed that some of the qualitative features of the experiment can be reproduced theoretically using a classical model similar to that in[20], which has been extended to two dimensions.

Up to now, a wide variety of BMs have been theoretically investigated and also demonstrated for various systems [3][4]. The main features of our BM can be qualitatively described using a purely classical model shown in section 4. Nevertheless, the coupling between the potentials is driven by quantum jumps as resulting from spontaneous emission, which essentially is a quantum mechanical feature. Together with quantised motion, this may open the way for creation of a quantum Brownian motor. Due to the generality of our scheme, applicability to chemical and/or biological systems may also be possible.

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