Characterization of stationary states in random walks with stochastic resetting

Vicenç Méndez and Daniel Campos
Grup de Física Estadística. Departament de Física. Facultat de Ciències. Edifici Cc. Universitat Autònoma de Barcelona, 08193 Bellaterra (Barcelona) Spain

It is known that introducing a stochastic resetting in a random-walk process can lead to the emergence of a stationary state. Here we study this point from a general perspective through the derivation and analysis of mesoscopic (Continuous-Time Random Walk) equations for both jump and velocity models with stochastic resetting. In the case of jump models it is shown that *stationary states* emerge for any shape of the waiting-time and jump length distributions. The existence of such state entails the saturation of the mean square displacement to an universal value that depends on the second moment of the jump distribution and the resetting probability. The transient dynamics towards this saturation and towards the stationary state depends on how the waiting time probability density function decays with time. If the moments of the jump distribution are finite then the tail of the stationary distributions is universally exponential, but for Lévy flights these tails decay as a power-law whose exponent coincides with that from the jump distribution. For velocity models we observe that the stationary state emerges only if the distribution of flight durations has finite moments of lower-order; otherwise, as occurs for Lévy walks, the stationary state does not exist and the mean square displacement grows ballistically or superdiffusively, depending on the specific shape of the distribution of movement durations.

PACS numbers:

I. INTRODUCTION

Random walks represent a recurrent tool to explore transport in systems subject to noise, fluctuations and/or uncertainty. In many applications, such walks can be interrupted (either by the moving particle/individual itself or by external forces) in such a way that the walker is brought back to its initial position and allowed to continue its movement from there newly. If this process is also driven by some noisy/fluctuating force we term it stochastic resetting. For example, in many spatial searches a natural tendency of living organisms is to return to the origin and start the search again after an unsuccessful excursion [1]. This is meaningful for example in foraging or other movement processes in animals which are often constrained by the presence of predators or other threats that can lead to the interruption of the movement as a risk-averse strategy or as a form of sheltering [2]. In a different context, stochastic resetting can be useful to describe information-spreading [3] or searches through graphs [4], particularly in Internet or other communication networks. In most of these examples, the focus is put in understanding the effects of stochastic resetting as a mechanism to enhance search efficiency (measured as the mean first passage time to a given target) under uninformed (random) search scenarios; an idea which has been explored in a formal way both for the Brownian motion case [5] and for exponential [6] and Lévy flights [7].

Another interest of resetting is in the dramatic effects it has over the stationary properties of transport processes. In [5, 8, 9] the authors found how Brownian particles subject to stochastic resetting evolve towards a nonequilibrium stationary state different from a Gaussian distribution due to the non-vanishing flux introduced by the resetting, which violates the detailed balance condition. However, up to date few attention has been put in studying the general conditions under which this nonequilibrium stationary state is expected to emerge. Moreover, some additional magnitudes which are typically of interest in random-walk processes, as the mean square displacement (MSD) (whose behavior can also give significant hints about the dynamics in the stationary regime), have not been computed for these situations. In this work we try to fill this gap by proposing a mesoscopic framework (based on the Continuous-Time Random Walk scheme) which includes both jump and velocity models. It is confirmed in general that the presence of the stationary state implies that the MSD grows monotonically towards a saturating value that equals the second moment of the distribution at the stationary state. If such saturating value does not exist, this is a signature that the stationary state is never reached (we will see this is the case for Lévy walks and similar processes). The relaxation dynamics towards this value depends on the corresponding waiting-time distribution. We consider several different situations with Markovian and non-Markovian distributions in both jump and velocity models. Specifically we review and extend here the results for the Lévy flight case (which has already been discussed in [5, 7]) and we study for the first time the case of Lévy walks as a particular case of a velocity model.

II. RESETTING-CTRW

We consider a unidimensional random-walk process starting from $x=0$ such that the walker has the possibility to reset its position to that origin whenever a single displacement is completed. We will denote by $X_1, X_2, ...$

the successive positions of the particle after the first, second,... event, where each event can be either a displacement or a reset. If we define r as the resetting probability, then the position X_{i+1} of the particle after the $(i+1)$-th event is chosen as $X_{i+1} = 0$ with probability r, provided the previous (i-th) event was not a reset. Otherwise, we choose $X_{i+1} = X_i + Z_i$, where the displacement length Z_i is a random variable drawn from the probability distribution function (pdf) $\Phi(x)$. We term this resetting mechanism as being *subordinated to displacements* [6] since the statistics of the displacements determines in part the rate at which resetting will occur. Note that this mechanism is slightly different to that proposed in [7] since here we explicitly require that a reset cannot occur immediately after another one. This choice has the advantage that the dynamics in the limit $r \to 1$ will consist of successive 'one-displacement' excursions separated by reset events. Otherwise, if consecutive resets were allowed then for $r \to 1$ the particle would remain permanently at $x = 0$, which for some applications may be unrealistic.

In order to take into account how different motion patterns affect the emergence of a stationary state in this system, we compare here jump and velocity models. In the former, walkers "movements" consist of instantaneous jumps which are separated by random waiting-times or pauses between them. In the latter these events consist of displacements done with finite speed v_0, so the distance travelled during one of these events and the movement duration are coupled variables.

Denote by $j(x,t)$ the density of particles starting a displacement from x at time t if it was initially at $x = 0$. A mesoscopic balance equation for $j(x,t)$ can be written then as

$$\begin{aligned} j(x,t) &= \delta(t)\delta(x) \\ &+ (1-r)\int_0^t dt' \int_{-\infty}^{\infty} dx' \Psi(x',t') j(x-x',t-t') \\ &+ r\delta(x)\int_0^t dt' \int_{-\infty}^{\infty} dx \varphi(t') j(x,t-t'). \end{aligned} \quad (1)$$

Here, $\Psi(x,t)$ is the joint probability of performing a displacement of length x during time t and the pdf for waiting times (in the jump's model) or movement duration (in the velocity's model) is given by $\varphi(t) = \int \Psi(x,t)dx$. So, the last term in (1) implicitly indicates that the resetting can only occur after a movement event has been completed.

The density of particles located at x at time t is given by

$$P(x,t) = \int_0^t dt' \int_{-\infty}^{\infty} dx' \phi(x',t') j(x-x',t-t'). \quad (2)$$

where $\phi(x,t)$ is just the probability that a single displacement has not finished yet after having travelled during a time t and having covered (either to left or right) a distance x. We define the Fourier-Laplace transform of an arbitrary function $g(x,t)$ as

$$g(k,s) = \int_0^\infty dt e^{-st} \int_{-\infty}^{\infty} dx e^{-ikx} g(x,t). \quad (3)$$

Transforming both equations (1) and (2) to the Fourier-Laplace space and combining them we obtain

$$P(k,s) = \left(1 + r\frac{\varphi(s)}{1-\varphi(s)}\right) \frac{\phi(k,s)}{1-(1-r)\Psi(k,s)}. \quad (4)$$

Another quantity of interest in this work is the flux density of particles $j(x,t)$, which in the Fourier-Laplace space is given by

$$j(k,s) = \left(1 + r\frac{\varphi(s)}{1-\varphi(s)}\right) \frac{1}{1-(1-r)\Psi(k,s)}. \quad (5)$$

III. JUMP MODEL

When displacement distances and waiting times are considered uncoupled random variables we can write $\phi(x,t) = \Phi(x)\varphi^*(t)$ and $\Psi(x,t) = \Phi(x)\varphi(t)$, where $\varphi^*(t) = \int_t^\infty dt'\varphi(t')$ is the survival probability of $\varphi(t)$, i.e., the probability for the particle not to jump away until time t. In this particular situation, Eq. (4) becomes

$$P(k,s) = \frac{1}{s}\frac{1-(1-r)\varphi(s)}{\Phi(k)^{-1}-(1-r)\varphi(s)}, \quad (6)$$

and from Eq. (5) the flux density of particles takes the form

$$j(k,s) = \left(\frac{1}{\varphi(s)} + \frac{r}{1-\varphi(s)}\right)\frac{1}{\varphi(s)^{-1}-(1-r)\Phi(k)}. \quad (7)$$

A. Relaxation towards the stationary state

In the large time limit $t \to \infty$ (which is equivalent to the limit $s \to 0$ in the Laplace space) a waiting time pdf with finite first moment $\langle t \rangle$ (mean waiting time) can be expanded through $\varphi(s) \simeq 1 - s\langle t \rangle$, and so Eq. (6) reads

$$P(k,s) \simeq \frac{1}{s}\frac{[r+(1-r)s\langle t\rangle]\Phi(k)}{1+s\langle t\rangle-(1-r)\Phi(k)}. \quad (8)$$

This can be inverted back to the real space in time, after some algebra, to obtain

$$P(k,t) \simeq P_s(k)\left[1 + \frac{1-\Phi(k)}{r\Phi(k)} e^{-\frac{1-(1-r)\Phi(k)}{(1-r)\Phi(k)}\frac{t}{\langle t\rangle}}\right]. \quad (9)$$

This reflects an exponential relaxation towards the stationary state

$$P_s(k) = \frac{r\Phi(k)}{1-(1-r)\Phi(k)} \quad (10)$$

since $0 < \Phi(k) \leq 1$ in the Fourier space, or towards

$$P_s(x) = \frac{r}{\pi} \int_0^\infty \frac{\cos(kx)dk}{\Phi(k)^{-1} - 1 + r} \tag{11}$$

in the real space. This is the stationary state distribution for any symmetric jump pdf $\Phi(x)$. This is a non-equilibrium steady state that is formed and sustained as a result of the permanent influx of particles to the origin due to the resetting process.

Let us deal now with the case where $\varphi(t)$ lacks finite moments. Consider the power-law pdf for waiting times that in the Laplace space reads $\varphi(s) = [1+(s\tau)^\gamma]^{-1}$ with $0 < \gamma < 1$. Inserting this into Eq. (6) we get

$$P(k,s) = \frac{1}{s} \frac{[r + (s\tau)^\gamma]\Phi(k)}{1 + (s\tau)^\gamma - \Phi(k)(1-r)}. \tag{12}$$

This equation can be inverted exactly by Laplace as follows

$$\begin{aligned} P(k,t) &= \frac{r\Phi(k)}{\tau^\gamma} t^\gamma E_{\gamma,\gamma+1}\left(-a(k)\frac{t^\gamma}{\tau^\gamma}\right) \\ &+ \Phi(k) E_\gamma \left(-a(k)\frac{t^\gamma}{\tau^\gamma}\right) \end{aligned} \tag{13}$$

where $E_\gamma(z)$ and $E_{\gamma,\gamma+1}(z)$ are Mittag-Leffler and Generalized Mittag-Leffler functions, and $a(k) = 1 - (1 - r)\Phi(k)$. Making use of the asymptotic expansions

$$\begin{aligned} E_\gamma(z) &= -\frac{1}{z\Gamma(1-\gamma)} + O(z^{-2}), \\ E_{\gamma,\gamma+1}(z) &= \frac{E_\gamma(z) - 1}{z} = -z^{-1} + O(z^{-2}) \end{aligned} \tag{14}$$

as $|z| \to \infty$ if $0 < \gamma < 1$, we can expand Eq. (13) for $t \to \infty$ to get

$$P(k,t) \simeq r\frac{\Phi(k)}{a(k)} \left[1 + \frac{1}{r}\left(\frac{\tau}{t}\right)^\gamma + ...\right]. \tag{15}$$

So that, the relaxation towards the stationary distribution follows a power law decay.

B. Properties of the stationary state

The shape of the stationary state given by (11) can only be found exactly for some particular cases. For example, consider the exponential jump pdf

$$\Phi(x) = \frac{1}{2\lambda} \exp(-|x|/\lambda) \tag{16}$$

with $\lambda > 0$ (also known as Laplace kernel). Its Fourier transform reads $\Phi(k) = (1 + k^2\lambda^2)^{-1}$. Then, from (11) we have that the stationary state reads

$$P_s(x) = \frac{\sqrt{r}}{2\lambda} e^{-|x|\frac{\sqrt{r}}{\lambda}},$$

an expression that has been obtained also previously in [5].

Let us inspect now the form of the stationary distribution from a more general perspective. First, we consider the diffusive limit, i.e. expand the jumps pdf $\Phi(x)$ up to the second moment $\Phi(k) \simeq 1 - \sigma^2 k^2/2$ as $\sigma k \ll 1$. In this case Eq. (11) reduces to

$$P_s(x) \simeq \frac{\sqrt{2r}}{2\sigma} e^{-|x|\frac{\sqrt{2r}}{\sigma}}, \text{ for } |x| \gg \sigma. \tag{17}$$

Therefore, all jump distance pdf's with finite moments have a stationary state decaying as in (17), where σ^2 is its second moment, $\sigma^2 = \int x^2 \Phi(x) dx$. It has been found recently [9] how such a stationary state is transiently followed in the case of Brownian motion by a still further region (for $|x| \leq \sqrt{4Drt}$, with D the diffusion coefficient of the Brownian walker) which includes those particles that have not experienced yet the effect of resetting, and that becomes eventually negligible in the large-time regime.

Besides the asymptotic behaviour, note that the shape of the stationary state close to $x = 0$ can also be approximated by expanding the $\cos(kx)$ in power series in Eq. (11). By doing that we have

$$P_s(x) \simeq a_0 - a_1 x^2, \quad \text{for } x \to 0 \tag{18}$$

where

$$\begin{aligned} a_0 &= \frac{r}{\pi} \int_0^\infty \frac{dk}{\Phi(k)^{-1} - 1 + r} \\ a_1 &= \frac{r}{2\pi} \int_0^\infty \frac{k^2 dk}{\Phi(k)^{-1} - 1 + r} \end{aligned} \tag{19}$$

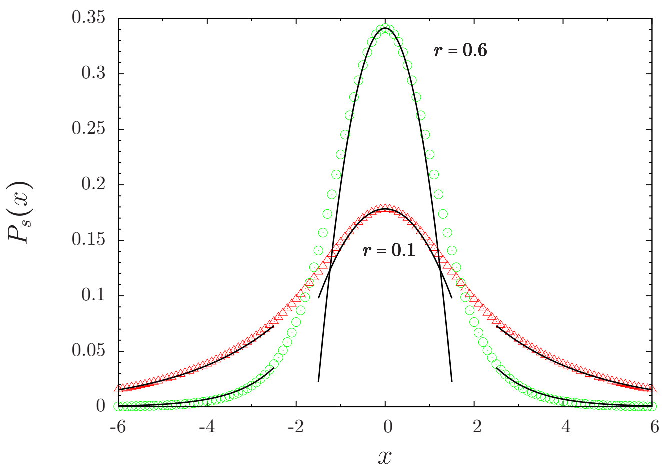

FIG. 1: (Color online) $P_s(x)$ for the Gaussian jump pdf for different values of the probability of resetting r. The results calculated from Eq. (11) have been drawn with symbols $\sigma = 1$. The solid curves indicate the approximations to the tails and the central part of $P_s(x)$.

In Figure 1 we plot the stationary state reached when the jump pdf obeys a Gaussian distribution $\Phi(x) =$

$[\sigma\sqrt{2\pi}]^{-1}\exp(-x^2/2\sigma^2)$. It is seen how the stationary distribution is more peaked around $x=0$ as r tends to 1, i.e., as the probability of reseting increases. The lines in the tails correspond to the exponential decay predicted by Eq. (17). Likewise, the lines of the central part correspond to the approximated solution given in Eq. (18).

We consider now the case of a Lévy distribution for jumps for which $\Phi(k)=e^{-\sigma^\mu|k|^\mu}$, with $1\le\mu\le 2$. Since $0<\Phi(k)<1$ then we can convert the right hand side of (10) into the sum

$$\begin{aligned} P_s(k) &= \frac{r}{1-r}\sum_{j=1}^{\infty}(1-r)^j\Phi(k)^j \\ &= \frac{r}{1-r}\sum_{j=1}^{\infty}(1-r)^j e^{-j\sigma^\mu|k|^\mu} \\ &= \frac{r}{\sigma(1-r)}\sum_{j=1}^{\infty}(1-r)^j j^{-1/\mu}L_\mu\left[\frac{|x|}{\sigma j^{1/\mu}}\right], \end{aligned} \quad (20)$$

where $L_\mu[x]$ is the normalized Lévy density defined by

$$L_\mu\left[\frac{|x|}{\sigma}\right]=\sigma\int_{-\infty}^{\infty}e^{-ikx-\sigma^\mu|k|^\mu}dx.$$

Expression (20) can be written in terms of Fox functions $H^{p,q}_{n,m}(z)$ [10] in the form

$$\begin{aligned} P_s(x) &= \frac{r\pi}{\mu(1-r)|x|} \\ &\times \sum_{j=1}^{\infty}(1-r)^j H^{1,1}_{2,2}\left[\frac{|x|}{\sigma}j^{\frac{-1}{\mu}}\,\middle|\,\begin{matrix}(1,\frac{1}{\mu}), & (1,\frac{1}{2})\\ (1,1), & (1,\frac{1}{2})\end{matrix}\right] \\ &= \frac{r\mu}{(1-r)|x|}\sum_{n=1}^{\infty}a(r,n,\mu)\left(\frac{\sigma}{|x|}\right)^{n\mu} \end{aligned} \quad (21)$$

for $|x|\gg\sigma$, where

$$a(r,n,\mu)\equiv\frac{(-1)^{n+1}\Gamma(n\mu)}{(n-1)!}\sin\left(\frac{n\mu\pi}{2}\right)\sum_{j=1}^{\infty}j^n(1-r)^j.$$

The tail of this stationary distribution behaves as

$$P_s(x)\simeq\frac{\Gamma(\mu)\sin(\pi\mu/2)}{r}\frac{\sigma^\mu}{|x|^{\mu+1}}, \quad (22)$$

just as the Lévy distribution in the real space. So, unlike the case where the jump pdf has finite moments the tail decays as that of the stationary distribution. In Figure 2 we repeat the same analysis as for Figure 1 but now for a Cauchy distribution

$$\Phi(x)=\frac{a/\pi}{x^2+a^2},$$

which corresponds to a Lévy distribution with $\mu=1$ and $\sigma=a$. The lines showed in the Figure 2 correspond to the approximations for small x prescribed in Eq. (18). Again, the stationary distribution is more peaked around $x=0$ for higher values of r and the tails decay heavily. To study more accurately the behaviour of the tails we show in Figure 2 (right) the tails for the same curves in a log-log scale.

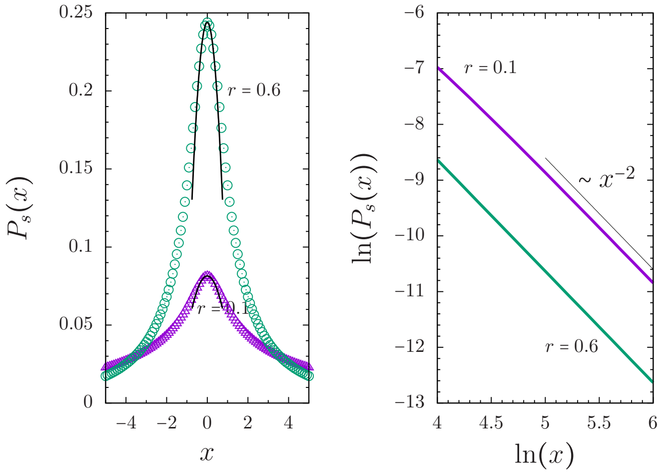

FIG. 2: (Color online) Left panel: $P_s(x)$ (symbols) for the Cauchy jump pdf for different values of the probability of resetting r. The symbols correspond to the solutions calculated from Eq. (11) taking $a=1$. Right panel: Log-log plot for $P_s(x)$ for the Cauchy jump pdf for different values of the probability of resetting r. The solid curves have been calculated from Eq. (11) taking $a=1$.

The flux density $j(x,t)$ also decays towards a stationary value $j_s(x)$ in the limit $t\to\infty$ but when the waiting time pdf has finite moments only. From Eq. (5), and proceeding in the same way as we have done for $P(x,t)$, the stationary flux density reads

$$j_s(x)=\frac{r}{2\pi\langle t\rangle}\int_{-\infty}^{\infty}\frac{e^{ikx}dk}{1-(1-r)\Phi(k)}. \quad (23)$$

If $\Phi(x)$ is exponentially distributed as in Eq. (16) then from (23) we find the exact solution

$$j_s(x)=\frac{(1-r)\sqrt{r}}{2\langle t\rangle\lambda}e^{-\frac{|x|}{\lambda}\sqrt{r}}+\frac{r}{\langle t\rangle}\delta(x). \quad (24)$$

Additionally, we consider again the case of a Lévy distribution for the jump pdf to compute the flux density at the stationary state. The trick is to convert the integrand of Eq. (23) into a sum and then take the inverse Fourier transform to end up with

$$j_s(x)=\frac{r}{\sigma\langle t\rangle}\sum_{j=0}^{\infty}\frac{(1-r)^j}{j^{1/\mu}}L_\mu\left[\frac{|x|}{\sigma j^{1/\mu}}\right] \quad (25)$$

Note that in all cases the flux density diverges at $x=0$ due to the effect of incoming particles from resetting, except trivially for $r=0$ (this is, when we remove resetting).

C. MSD

To compute the MSD we have to assume that $\Phi(x)$ has finite moments. By using Eq. (6) we have

$$\left\langle x^2(s)\right\rangle = -[\partial_{kk}P(k,s)]_{k=0} = \frac{\left\langle l^2\right\rangle}{s[1-(1-r)\varphi(s)]}, \quad (26)$$

where $\left\langle l^2\right\rangle$ is the second moment of $\Phi(x)$, i.e. $\left\langle l^2\right\rangle = [\Phi''(k)]_{k=0}$. If we consider the Markovian case where the waiting-time pdf is exponentially distributed, i.e. $\varphi(t) = \tau^{-1}e^{-t/\tau}$ (or, equivalently, $\varphi(s) = [1+s\tau]^{-1}$), then employing Eq. (26) it follows

$$\left\langle x^2(t)\right\rangle = \frac{\left\langle l^2\right\rangle}{r}\left[1-(1-r)e^{-rt/\tau}\right]. \quad (27)$$

This shows an exponential convergence to the asymptotic value $\left\langle l^2\right\rangle/r$ with a characteristic relaxation time τ/r.

Let us now consider the non-Markovian case where the waiting -time pdf decays as a power law in time. In the large time limit ($s\to 0$) we consider $\varphi(s)\simeq 1-(s\tau)^\gamma$. Then, (26) reads

$$\left\langle x^2(t)\right\rangle = \frac{\left\langle l^2\right\rangle}{(1-r)\tau^\gamma}t^\gamma E_{\gamma,\gamma+1}\left(-\frac{rt^\gamma}{(1-r)\tau^\gamma}\right). \quad (28)$$

Using the expansion in (14) into (28) leads to the result

$$\left\langle x^2(\infty)\right\rangle = \frac{\left\langle l^2\right\rangle}{r} \quad (29)$$

which is independent of the anomalous exponent. This result illustrates that the MSD converges to the above result when $t\to\infty$. On the other hand we can check that this result is precisely equal to the second moment of the pdf $P_s(x)$. By its definition, the second moment is

$$\left\langle x^2(\infty)\right\rangle = \int_{-\infty}^{\infty} x^2 P_s(x)dx = -[\partial_{kk}P(k,\infty)]_{k=0} = \frac{\left\langle l^2\right\rangle}{r}$$

where we have used (10). The expression (29) corroborates the result that if $\Phi(x)$ lacks finite moments then the same happens for the stationary distribution too. In Figure 3 we plot the MSD computed from Eq. (28). The relaxation to the asymptotic value given by Eq. (29) follows the power law $t^{-\gamma}$. As a result the relaxation is governed by the value of the exponent γ, as can be checked in Figure 3 for values from $\gamma = 0.8$ to $\gamma = 0.4$.

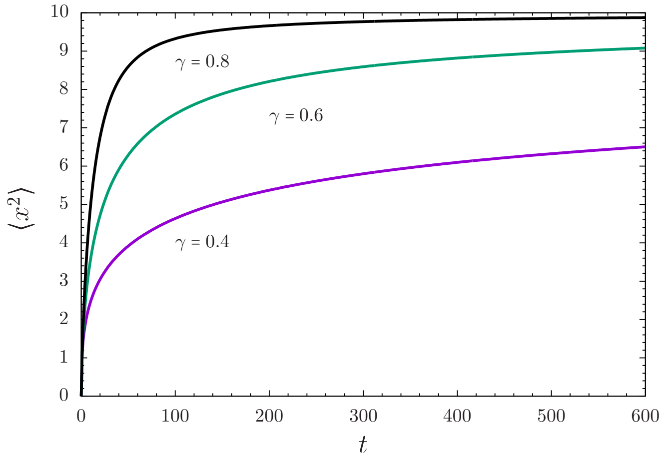

FIG. 3: (Color online) Plot of the MDS relaxing to the asymptotic value $\left\langle l^2\right\rangle/r$. In this case the jump pdf has finite moments but the waiting-time pdf is a power-law. Here $\tau = \left\langle l^2\right\rangle = 1$, $r = 0.1$.

IV. VELOCITY MODEL

A. Stationary state

Here $\Psi(x,t) = \Phi(x|t)\varphi(t)$ is the joint probability of performing a movement of length $|x| = v_0 t$ with constant speed v_0 during time t. So, the quantity $\Phi(x|t)$ is the conditional probability to perform a movement of duration t to the right or left with probability 1/2 and with length $v_0 t$. Therefore,

$$\begin{aligned}\Psi(x,t) &= \frac{1}{2}\left[\delta(x-v_0t)+\delta(x+v_0t)\right]\varphi(t)\\ &= \frac{1}{2v_0}\delta\left(t-\frac{|x|}{v_0}\right)\varphi(t) \quad (30)\end{aligned}$$

The probability of having completed a movement of distance x is $\Phi(x)$, and can be computed from (30)

$$\Phi(x) = \int_0^\infty dt\Psi(x,t) = \frac{1}{2v_0}\varphi\left(\frac{|x|}{v_0}\right). \quad (31)$$

On the other hand,

$$\phi(x,t) \equiv \frac{1}{2}\left[\delta(x-v_0t)+\delta(x+v_0t)\right]\varphi^*(t) \quad (32)$$

is just the probability that a single movement has not finished yet after having traveled during a time t and having covered (either to left or right) a distance $v_0 t$.

To compute the density at the steady state let us take the limit $s\to 0$ in Eq. (4)

$$\begin{aligned}P(k,s) &\simeq \frac{\phi(k,s=0)}{1-(1-r)\Psi(k,s=0)}\lim_{s\to 0}\left(1+r\frac{\varphi(s)}{1-\varphi(s)}\right)\\ &= F(k)\lim_{s\to 0}\left(1+r\frac{\varphi(s)}{1-\varphi(s)}\right), \quad (33)\end{aligned}$$

where we have defined

$$F(k) \equiv \frac{-\text{Im }\varphi(ikv_0)}{kv_0\left[1-(1-r)\text{Re }\varphi(ikv_0)\right]} \quad (34)$$

Consider the case where the pdf of movement durations has finite moments. In the large time limit we can use the expansion $\varphi(s)\simeq 1-s\langle t\rangle$, so in this case

$$\lim_{s\to 0}\left(1+r\frac{\varphi(s)}{1-\varphi(s)}\right) \simeq \frac{r}{\langle t\rangle s}+\ldots \quad (35)$$

Inverting by Fourier (33) and taking into account (34) and (35), the general expression for the stationary state reads

$$P_s(x) = \frac{r}{\pi\langle t\rangle}\int_0^\infty \cos(kx)F(k)dk \tag{36}$$

If we want to investigate the tail of the stationary distribution we have to take the large time limit which in this case is equivalent to the large space limit, since both variables are coupled through v_0. Considering that the pdf of movement durations has finite moments, in the large time limit it can be written as $\varphi(s) \simeq 1 - s\langle t\rangle + s^2\langle t^2\rangle/2$. After inserting this into (34) we get

$$F(k) = \frac{1}{r + \frac{1}{2}(1-r)\langle t^2\rangle k^2 v_0^2}.$$

Introducing this result into (36) we finally obtain

$$P_s(x) \simeq \frac{\sqrt{r}}{\sqrt{2(1-r)\langle t^2\rangle}v_0}\exp\left(-\frac{\sqrt{2r}}{\sqrt{(1-r)\langle t^2\rangle}v_0}|x|\right). \tag{37}$$

This is the stationary distribution for a velocity model when $|x|$ is large. Unlike the jump model, in the velocity model the $P_s(x)$ depends also on the second moment of the movement duration pdf due to the time-space coupling.

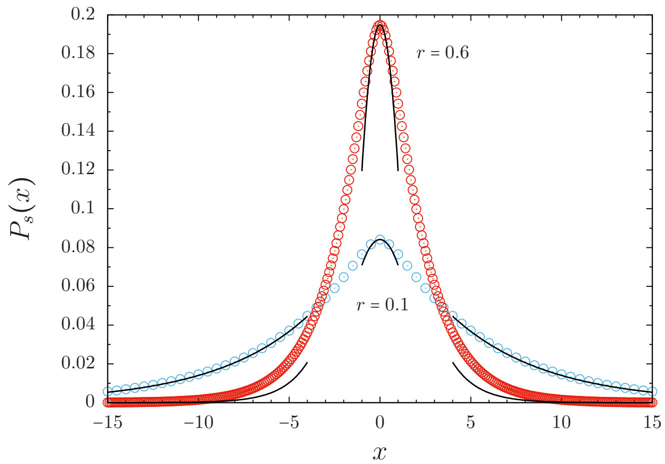

FIG. 4: (Color online) Stationary distribution (symbols) for a pdf movement durations $\varphi(t) = te^{-t/\tau}/\tau^2$ for different values of r. $v_0 = \tau = 1$.

In Figure 4 we have plotted the exact results given by the Eq. (36) when we make use of the Gamma distribution $\varphi(t) = te^{-t/\tau}/\tau^2$. The lines in the tails correspond to the theoretical asymptotic result predicted by Eq. (37), and the lines at the central part of the distributions are computed through $P_s(x) \simeq a_2 - a_3x^2$, where

$$a_2 = \frac{r}{\pi\langle t\rangle}\int_0^\infty F(k)dk$$

$$a_3 = \frac{r}{\pi\langle t\rangle}\int_0^\infty k^2F(k)dk.$$

This is the highest-order possible approximation to the central part since the integrals

$$\int_0^\infty k^nF(k)dk$$

diverge for $n > 2$.

Consider now a non-Markovian example with a power-law pdf of movement durations $\varphi(s) = [1 + (s\tau)^\gamma]^{-1}$. In this case

$$\lim_{s\to 0}\left(1 + r\frac{\varphi(s)}{1-\varphi(s)}\right) \simeq \frac{r}{(s\tau)^\gamma} + ...$$

After simplifying (33) we find

$$P(k,s) \simeq \frac{1}{(s\tau)^\gamma}\frac{(k\tau v_0)^{\gamma-1}\sin(\pi\gamma/2)}{1 + \frac{1+r}{r}\cos(\pi\gamma/2)(k\tau v_0)^\gamma + r^{-1}(k\tau v_0)^{2\gamma}}.$$

By inverting the Laplace transform, the factor $(s\tau)^{-\gamma}$ turns into a factor $t^{\gamma-1}$ which tends to 0 as $t \to \infty$ since $0 < \gamma < 1$. Then, when the pdf of movement durations lacks finite moments there is no stationary state and the relaxation to the stationary state $P_s(x) = 0$ follows the power-law decay $t^{\gamma-1}$.

B. MSD

We finally explore the behaviour of the MSD when the pdf of movement durations and the pdf of movement distances have or does not have finite moments. Starting from the general expression (4) we find after some calculations

$$\begin{aligned}\left\langle x^2(s)\right\rangle &= -[\partial_{kk}P(k,s)]_{k=0} \\ &= v_0^2\left[\frac{(\varphi^*(s))''}{1-\varphi(s)} + \frac{1-r}{s}\frac{\varphi''(s)}{1-(1-r)\varphi(s)}\right]\end{aligned} \tag{38}$$

where the symbol $''$ means second derivative respect to s. For the case of movement duration pdf's with finite moments we consider the asymptotic expansion $\varphi(s) \simeq 1 - s\langle t\rangle + s^2\langle t^2\rangle/2 - s^3\langle t^3\rangle/6$. Hence, in the large time limit Eq. (38) reduces to

$$\left\langle x^2(\infty)\right\rangle = v_0^2\left(\frac{\langle t^3\rangle}{3\langle t\rangle} + \frac{1-r}{r}\left\langle t^2\right\rangle\right). \tag{39}$$

For the case of the exponential distribution $\varphi(t) = e^{-t/\tau}/\tau$ we have seen that there exists a stationary state so it is expected that the MSD tends to a constant. To find the time dependence of the MSD we make use of (38) and get

$$\left\langle x^2(t)\right\rangle = \frac{2v_0^2\tau^2}{r}\left(1 + \frac{re^{-t/\tau} - e^{-rt/\tau}}{1-r}\right) \tag{40}$$

Performing the limit $t \to \infty$ we find

$$\langle x^2(\infty) \rangle = \frac{2v_0^2\tau^2}{r}.$$

This result equals the second moment of the pdf at the stationary state as already happened in the jump model. Effectively, from (34) we have

$$F(k) = \frac{\tau}{r + k^2 v_0^2 \tau^2}$$

and together with (36)

$$\begin{aligned} \langle x^2(\infty) \rangle &= \int_{-\infty}^{\infty} x^2 P_s(x) dx \\ &= \frac{\sqrt{r}}{v_0\tau} \int_0^{\infty} x^2 e^{-x\frac{\sqrt{r}}{v_0\tau}} dx = \frac{2v_0^2\tau^2}{r}. \end{aligned} \quad (41)$$

Finally, let us consider the case of Lévy walks. They are random walks where the displacements are performed with finite velocity but the jump distribution or the movement duration pdf (lined through Eq. (31)) have tails which decay according to power-laws, so higher-order moments are lacking. For example the pdf $\varphi(s) = [1 + (s\tau)^\gamma]^{-1}$, with $0 \leq \gamma \leq 1$, lacks all moments including its mean value. If we insert this expression into Eq. (38) we find

$$\langle x^2(t) \rangle = \frac{(2-\gamma)(1-\gamma)}{2} v_0^2 t^2 \quad \text{as} \quad t \to \infty$$

which corresponds to a ballistic transport regime. This behavior is due to the fact that the probability density of particles $P(x,t)$ is a sandwich between two ballistic peaks located at $x = \pm v_0 t$. Another possibility is to consider $\varphi(s) \simeq 1 - s\langle t \rangle + As^\mu$ with $1 < \mu < 2$. In this case there exists the first moment (but not the second-order or higher). Calculation of the MSD for this case gives from Eq. (38) a superdiffusive behaviour

$$\langle x^2(t) \rangle = \frac{(2-\mu)(1-\mu)}{\Gamma(4-\mu)\langle t \rangle} v_0^2 A t^{3-\mu} \quad \text{as} \quad t \to \infty.$$

Again, the result does not depend on the probability of resetting explicitly because the asymptotic behaviour is dominated by the fraction of particles which have not experienced resetting yet. So, in both cases studied (and in general for Lévy walks) we find that there is no stationary state, contrary to what happened for the case of Lévy flights.

V. SUMMARY

In this work we have studied the conditions for the existence of a resetting-induced stationary state and so a saturation value for the MSD in the case of a jump model where jumps distances and waiting times are independent random variables and the walker is submitted to stochastic resetting with probability r. The tail of the stationary distribution is exponential if the jump pdf has finite moments but when the jumps pdf decays as a power law the stationary pdf also decays as a power-law with the same exponent. The MSD grows exhibiting an exponential saturation if the waiting time pdf has finite moments or saturates as a power law in time if the waiting time pdf also decays as a power-law. The saturating value is always $\langle l^2 \rangle / r$, where $\langle l^2 \rangle$ is the second moment of the jump pdf. In consequence, the finiteness of the second moment (at least for isotropic motion) determines the saturation of the MSD. The situation under advective or biased movement will require further examination and will be the focus of a forthcoming work.

In the velocity model, where the movement of the walker is performed at constant speed v_0, the movement duration pdf defined univocally the jump displacement pdf. When the movement duration pdf has finite moments there is a stationary state with again an exponential tail and the MSD saturates always to the value $2v_0^2\tau^2/r$. When it lacks first or second order moments, as for the Lévy walks case, there is no stationary state and the MDS grows ballistically or superdiffusively, respectively. This is a consequence of our choice of resetting *subordinated to displacements*; we note that implementing resetting as an independent process of motion (as was done, for example, in [11] or [6]) then the stationary state will emerge (both for jump and velocity models) whenever the mean time for resetting is finite.

Acknowledgements.

This research has been partially supported by Grants No. FIS 2012-32334 and SGR 2013-00923. VM also thanks the Isaac Newton Institute for Mathematical Sciences, Cambridge, for support and hospitality during the CGP programme, where part of this work was undertaken.

[1] D. Campos, F. Bartumeus, V. Méndez and X. Espadaler. J. Roy. Soc. Interf. 11, 20130859 (2013).

[2] O. Bénichou, C. Loverdo, M. Moreau and R. Voituriez, Rev. Mod. Phys. **83**, 81 (2011).

[3] D. Liben-Nowell and J. Kleinberg, Proc. Natl. Acad. Sci. USA **105**, 4633 (2008).

[4] E. Gelenbe. Phys. Rev. E 82, 061112 (2010).

[5] M. R. Evans and S. N. Majumdar, Phys. Rev. Lett. **106**,

160601 (2011).
[6] D. Campos and V. Méndez. arXiv 1509.05641.
[7] L. Kusmierz, S.N. Majumdar, S. Sabhapandit and G. Schehr. Phys. Rev. Lett. **113**, 220602 (2014).220602 (2014).
[8] M.R. Evans and S.N. Majumdar. J. Phys. A: Math. Theor. 47, 285001 (2014).
[9] S.N. Majumdar, S. Sabhapandit and G. Schehr. Phys. Rev. E 91, 052131 (2015).
[10] B. J. West, et. al. Phys. Rev. E **55**, 99 (1997).
[11] M. Montero and J. Villarroel. Phys. Rev. E 87, 012116 (2013).