Surface dispersion in the Gulf of California

L. Zavala Sansón

Department of Physical Oceanography, CICESE, Mexico

A B S T R A C T

Surface dispersion is measured in the Gulf of California by means of Argos drifters released along this semi-enclosed, elongated basin. First, basic one-particle statistics (Lagrangian scales, absolute dispersion and diffusion coefficients) are estimated along and across the Gulf. Absolute dispersion shows a nearly ballistic regime during the Lagrangian time scale (<2 days) in both directions (it grows as \( \sim t^2 \), where \( t \) is time). During the subsequent 30 days, absolute dispersion enters a random-walk regime (\( \sim t \)) along the Gulf, while being saturated across the basin due to the lateral boundaries. Secondly, the analysis is extended to two-particle statistics (relative dispersion between pairs of drifters and Finite Scale Lyapunov Exponents, FSLE). Relative dispersion is nearly exponential in both directions during the first few days, though evidence is not conclusive. During the subsequent 30 days, it grows as \( \sim t^{1.5} \) along the Gulf, while being saturated across the basin again. It is shown that relative dispersion along the Gulf is proportional to \( t^4 \), where \( t \) represents a shifted time that depends on the initial separation of the particles. This form of the Richardson regime is consistently measured for particles that are sufficiently separated (30 km or more). The Richardson regime is verified with the FSLE for particle separations ranging from 30 to 140 km, approximately. The obtained dispersion properties are discussed in terms of the main circulation features within the basin, such as mesoscale vortices that occupy the width of the Gulf. These structures might retain buoys during days or weeks, thus preventing or delaying further displacements and therefore affecting the particle dispersion. The vortices are also an important mechanism to translate particles across the Gulf, between the Peninsula and the continent, thus promoting the saturation of dispersion along this direction.

© 2015 Elsevier Ltd. All rights reserved.

1. Introduction

The mechanisms governing the dispersion of a set of tracers immersed in a turbulent flow are of fundamental importance to understand and/or predict the fate of physical, chemical and biological properties in the oceans and in the atmosphere. This is usually a complex problem given the wide range of spatial scales involved in a turbulent geophysical flow, which might range from a few centimeters to thousands of kilometers. One of the methods to study oceanic dispersion consists of releasing a number of buoys (or balloons in the atmosphere) over the surface or at different depths (heights), and then calculating the statistical properties of their distribution as they spread. There is a wide variety of analytical tools and physical models for analysing the statistical behavior of Lagrangian tracers, some of which will be used in this study. For a recent review see LaCasce (2008) or Salazar and Collins (2009).

Here we present surface dispersion properties in the Gulf of California (hereafter denoted as GC) measured by means of surface drifters released all over the basin in different periods of the year. One and two-particle statistics shall be discussed. First, within the framework of single-particle analysis, we measure the time and space Lagrangian scales of the turbulent ambient flow, which are based on the velocity autocorrelation functions for each drifter (Taylor, 1921). Afterwards, we calculate the absolute dispersion, which indicates the representative (squared) separation of particles with respect to their initial position. Absolute dispersion is an appropriate measure of how far individual particles move away from their original position, while absolute diffusion indicates how fast the tracers are dispersed.

Secondly, two-particle statistics are analyzed with pairs of drifters. In contrast to absolute dispersion, relative dispersion measures the separation of two particles or, equivalently, the spread of a cloud of passive tracers. Some of the earliest ideas developed to quantify relative dispersion were proposed by Richardson (1926). A modern way to characterize relative dispersion is by means of the so-called finite scale Lyapunov exponents (Artale et al., 1997; Aurell et al., 1997). This method allows the determination of the relative diffusivity in terms of the scale of separation between pairs of drifters. This is a fundamental point that must be considered, since it is well-known since the work of Richardson that the dispersion properties strongly depend on the length scale of the spot of passive tracers. In particular, we shall
present measurements that support the presence of the Richardson regime, in which the relative diffusivity is proportional to the mean particle separation elevated to the 4/3 power (the so-called 4/3s Richardson law).

There are two important features that underline the relevance of the dispersion problem within the GC. First, the Gulf is a semi-enclosed sea with an open boundary at the southern part, facing the Pacific Ocean (Fig. 1). In contrast, several studies on dispersion have been focused on turbulent flows without the influence of lateral walls or, in the oceanic case, for open ocean regimes (e.g., Ollitrault et al., 2005). Lateral boundaries exert a strong influence on the circulation of surface drifters by limiting their motion. It will be shown, indeed, that Lagrangian dispersion is influenced by the long sides of the Gulf: the Baja California Peninsula to the west and the Mexican continental side to the east. Some studies on closed or semi-enclosed domains have been carried out in the Adriatic Sea (Falco et al., 2000; Lacorata et al., 2001) or in idealized systems (Artale et al., 1997).

A second relevant aspect is the shape of the Gulf, which is a very elongated basin, about 150–200 km wide and 1400 km long, somewhat longer than the Adriatic. This morphology plays a fundamental role on the typical circulation, which in turn strongly influences the dispersion. The surface circulation in the GC is characterized by the presence of mesoscale gyres and along-shore jets (Lavín and Marinone, 2003). The circulation at the northern part is defined by a single gyre that reverses seasonally, being cyclonic during the Summer and anticyclonic during the Winter (Lavín et al., 1997; Beier and Ripa, 1999). At the central and southern parts, between 23° and 28°N, a series of alternating vortices with horizontal dimensions comparable with the width of the Gulf (<200 km), is commonly observed (e.g., Pegau et al., 2002). According to historical hydrographic data, however, there are no definite patterns in the position of the gyres, nor in the season at which they appear nor in their sense of rotation (Figuerola et al., 2003). Recently, Lavín et al. (2013, 2014) described with great detail the hydrodynamic characteristics of a train of alternating mesoscale eddies in the southern GC during the Summer. According to numerical simulations (Zamudio et al., 2008), the geostrophic eddies are presumably formed as an intense jet flowing northwestward along the eastern boundary becomes unstable when interacting with coastal irregularities. This enhanced coastal current has been recently documented in the observations of Lavín et al. (2014). The presence of the developed vortices might play a relevant role in the dispersion of tracers by trapping particles during days or weeks, or by providing an effective advection mechanism between the long sides of the GC.

There are several practical motivations for studying dispersion in the GC. For instance, the spreading of biological material is particularly relevant to understand, given the rich biodiversity and the presence of several important fisheries. Recent studies have examined the connectivity properties within the GC, that is, the degree of association between different areas with respect to the motion of passive tracers, with which larval dispersal is represented in numerical simulations (Marinone, 2012). Regarding numerical studies, the Lagrangian circulation in the GC has been studied with two and three dimensional simulations (Marinone, 2006; Marinone et al., 2011). More in general, some numerical Lagrangian schemes simulate the spreading of passive tracers in the ocean as a diffusive process (Lumpkin and Elipot, 2010), and usually adopt suitable values of horizontal and vertical diffusivities. These parameters are often rather artificial or difficult to justify, and therefore direct estimations are usually made by means of drifters. (see e.g. Qian et al., 2014; Zhurbas et al., 2014).

In Section 2 the data set of drifters is presented together with some typical examples of buoy trajectories, in order to illustrate relevant features of the circulation in the GC. One and two-particle statistics are presented in Sections 3 and 4, respectively. The results are discussed in Section 5, and the main conclusions are presented in Section 6.

2. Buoys

2.1. Data set

The buoy data were obtained during different oceanographic campaigns carried out by Dr. Miguel Lavín and his team from CICESE (Mexico). Several groups of surface drifters (SVP, Pacific Gyre) were released between 2004 and 2006 inside and at the entrance of the GC. Some deployments were made in oceanographic cruises and some others from a commercial ferry traveling from/to the city of La Paz (~24°N at the Baja California Peninsula) to Mazatlan (~23°N at the continental side of Mexico). The surface drifters had a drogue centered at a 15 m depth. The buoys were tracked by means of the satellite ARGOS system, with position accuracy of 150–1000 m (Lumpkin and Pazos, 2006). Using this information, together with satellite images, the surface circulation of the GC was described by Lavín et al. (2014) during a timespan of more than two years.

![Fig. 1](image-url) (a) Trajectories of 61 Argos drifters inside the Gulf of California. Initial positions are indicated with dots. Buoy trajectories leaving the Gulf (south of 23°N) are discarded. (b) Lifetime of the 61 drifters. The time origin is 6 June 2004. The mean lifetime is \( T_m = 45 \pm 28 \) days. The coordinate system is rotated an angle \( \pi/5 \) to express the displacements and velocity components across an along the GC.
We selected the drifters whose trajectories (complete or partial) lied inside the Gulf (north of 23°N), and were at least 10 days long. The final data set consists of 61 surface buoys. Given the number of tracers and the almost two-year period of measurements, the whole surface of the Gulf is sampled, as shown in Fig. 1a, where the superposition of all the selected buoy trajectories is presented. The portion of the trajectories leaving or entering the Gulf are discarded (note that some drifters leave the basin at 23°N). The buoy deployment is not uniform along the Gulf: an important group of drifters was released at the southern region, between 23° and 25°N. A second set of buoys was deployed at the central Gulf, around 28°. A third, less numerous set of drifters was released at the northern part, at 30°N. The lifetimes of each buoy are presented in Fig. 1b. The duration of the records ranges between 10 and 90 days, with a mean of 45 days and a standard deviation of 28 days. The deployment of the drifters was not designed to measure dispersion properties but to study mesoscale circulation features, specially at the southern and central regions (Lavín et al., 2013). Even though, it shall be shown that one and two-particle statistics can be estimated with the available buoy trajectories.

The position and the horizontal velocity of each buoy were interpolated (without filtering) to regular 6 h intervals (Lavín et al., 2014). Let \( u(t) \) be the zonal (meridional) velocity component. Given the geographical orientation of the GC, it is convenient to rotate the coordinate system an angle \( \pi/5 \), such that the velocity components transform as

\[
\begin{align*}
u_1 &= u \cos(\pi/5) + v \sin(\pi/5), \quad (1) \\
u_2 &= -u \sin(\pi/5) + v \cos(\pi/5), \quad (2)
\end{align*}
\]

where subscripts 1 and 2 indicate the components across and along the Gulf, respectively. The dispersion measurements will be given in terms of the rotated components.

### 2.2. Typical trajectories

Before presenting the dispersion and diffusion measurements, it is convenient to show some typical buoy trajectories that reveal important features of the circulation inside the GC. First, we illustrate some of these motions by means of individual drifters. As mentioned in the Introduction, mesoscale vortices with horizontal dimensions comparable with the width of the Gulf are commonly observed. Fig. 2a shows the trajectory of a buoy that was released in Summer at the southern part; the drifter traveled northward until being trapped by an anticyclonic structure at about 27°N during more than one month. Afterwards, the buoy escapes and moves further north. Fig. 2b presents the trajectory of a buoy that was trapped by the seasonal cyclonic gyre at the northern Gulf.

The drifters are not necessarily captured by mesoscale structures. In some cases, the buoys might be advected along the basin in either direction during relatively short times, as shown in Fig. 2c and d. The buoy in panel c was released in Winter at about 28°N, and traveled southward before leaving the Gulf after two months. Drifter in panel d was deployed in Summer and moved northward along the entire basin in 2.5 months. In both cases the drifters exhibit a meandering motion but almost no trapping occurred. Apparently, all the behaviors presented in Fig. 2 (vortex trapping or meandering motion along the basin) will play a role on the dispersion properties.

The analysis of drifter pairs reveals further details of the circulation. Fig. 3a shows two buoys that are released at a close distance from each other (48 km, see caption), and follow a nearly parallel trajectory dictated by the presence of mesoscale vortices at the central-south region of the GC. The drifter pair was released in Summer and moved northward during more than one month. A similar case is presented in Fig. 3b for buoys released in Winter, which start at about 27°N and follow a southward trajectory before being separated at 25°N. A closer examination reveals that one of the buoys performs a few small-scale turns at the beginning, and again one month later, at 25°N, where it is trapped inside a gyre while the other drifter continues moving further southward.

A contrasting case is presented in Fig. 3c, where another pair of buoys immediately drifted in opposite directions after deployment at the central Gulf. Particle pairs also illustrate the presence of different scales of motion simultaneously, as shown in Fig. 3d: one of the drifters is trapped inside a mesoscale anti-clockwise gyre whose width is similar to that of the basin, while the other buoy presents a similar motion while experiencing additional, smaller anticyclonic turns at the periphery. The size of these gyres indicates the possible presence of smaller vortices compared with the width of the GC. We shall not discuss in detail the origin of such motions, but we point out that they presumably play a role in the dispersion properties.

### 3. One-particle statistics

Absolute dispersion (or one-particle statistics) is based on the mean-squared separation of particles from their position at a reference time. Using \( N \) Lagrangian trajectories one can calculate absolute dispersion along an arbitrary direction as (Provenzale, 1999)

\[
\overline{A_i^2}(t) = \frac{1}{N} \sum_{k=1}^{N} [x_i^k(t) - x_i^k(t_0)]^2,
\]

where \( x_i^k \) is the position of particle \( k \) in the \( i \)-direction (in two dimensions \( i = 1, 2 \), \( t \) is time and \( t_0 \) is the time at which the trajectories are released. The overbar indicates ensemble average. Absolute diffusion is defined as the temporal evolution of the dispersion

\[
K_i(t) = \frac{1}{2} \frac{d}{dt} \overline{A_i^2}.
\]

and measures how fast the particles are dispersed.

There are several methods to calculate dispersion and diffusion properties according with pre-defined assumptions on the turbulent character of the flow. The basic theory of single-particle dispersion was derived by Taylor (1921), and it is based on the assumption of homogeneous and stationary turbulence without a mean flow. In Section 4.1 we estimate the so-called Lagrangian scales (for time, length and diffusivity) according with Taylor’s theory. They indicate the corresponding scales at which the drifter still contains information about its past trajectory (due to the fluid motion in which it is embedded). In Section 4.2 we directly measure absolute dispersion from (3). This definition applies for particles released at the same position in different realizations of the turbulent flow (Provenzale, 1999). In practice, however, drifters in the ocean are released at different places and times. The assumption of homogeneous turbulence makes compatible both situations.

### 3.1. Lagrangian scales

The Lagrangian scales are estimated from the velocities of individual drifters by assuming that these are advected by a homogeneously turbulent flow. For these calculations we follow the procedure used by Poulain and Niiler (1989). Let \( u_i \) the velocity components of particle \( k \) along its trajectory. We obtain the residual velocity by subtracting a Lagrangian average over the trajectory during a time \( T \):

\[
u_i'(t) = u_i(t) - \frac{1}{T} \int_0^T u_i(t) \, dt.
\]

An alternative method is to subtract an Eulerian mean flow to the particle velocities, which is usually more difficult to obtain, and it
is subject to a number of additional uncertainties, such as the size of
the geographical bin chosen for each velocity vector (LaCasce,
2008).

The ‘memory’ of the drifter is contained in the Lagrangian auto-
covariance function defined as
\[ R_{ii}(s) = \frac{1}{T} \int_0^T u_i(t) u_i(t + s) dt, \]
which measures the autocorrelation of each velocity component
along the particle trajectory (so \( s \) is a time lag between two consec-
tutive positions). The residual velocity variance is the autocovariance
function with zero lag
\[ R_{ii}(0) = \frac{1}{T} \int_0^T u_i^2(t) dt = R_{00}. \]
In general, the autocovariance also depends on the initial time and
position of the particle (Poulain and Niiler, 1989). However, the
assumption of homogeneity and stationarity of the flow field allows
one to ignore such dependence, which is equivalent to assume that
all the particles were released at the same place, at the same time.

The period \( T \) is of critical relevance: it has to be long enough to
consider drifters with long autocorrelations, but short enough to
obtain a representative Lagrangian mean velocity in (5). We choose
this time as the mean lifetime of the whole set of drifters
(\( T = T_m = 45 \) days), except when the buoy lifetime \( T_f \) is shorter;
in that case \( T = T_f \). The results are very similar when selecting
times between 30 and 90 days instead of \( T_m \) (90 days corresponds
to use \( T = T_f \) for all buoys, i.e. to calculate the Lagrangian average
over the entire trajectories).

The turbulence is characterized by a time-independent value of
the variance \( R_{00} \). The Lagrangian time scale \( T^i_k \) is defined as
\[ T^i_k = \frac{1}{R_{00}} \lim_{\tau \to \infty} \int_0^\tau R_i(\tau) d\tau, \]
which is interpreted as the period during which the particle ‘re-
members’ its path, that is, the timespan during which its velocity
is well autocorrelated. An upper bound of \( T^i_k \) is obtained by integrat-
ing the autocorrelation function up to the first zero crossing at \( t_Z \).

The Lagrangian length scale can be computed as
\[ L^k_i = (R_{00})^{1/2} T^i_k = \frac{1}{(R_{00})^{1/2}} \int_0^{t_Z} R_i(\tau) d\tau. \]
Using \( T^k_i \) and \( L^k_i \), the diffusivity scales as
\[ K^i_k = \frac{(L^k_i)^2}{T^k_i} = \frac{R_{00}(T^k_i)}{\int_0^{t_Z} R_i(\tau) d\tau}. \]
Recent studies have pointed out that the integration of the autocor-
relation function to the first zero crossing implies an overestimation
of the diffusivity, because this function usually displays a pro-
nounced positive lobe at early times (Klocker et al., 2012; Zhurbas
et al., 2014). These authors discuss the simultaneous decaying and
oscillatory behavior of the autocorrelation function, the latter being
related with the phase speed of eddies embedded in a mean shear
flow. We shall show, however, that the estimation of diffusivities
using the first zero-crossing criterion leads to very similar results.
as those obtained with direct measurements of absolute dispersion, presented in next subsection.

The mean Lagrangian scales are measured by averaging over the ensemble of 61 drifters:

\[
\langle T_i, L_i, K_i \rangle = \frac{1}{N} \sum (T_k^i, L_k^i, K_k^i) \tag{11}
\]

The calculated values are shown in Table 1 (upper part). Note that the Lagrangian scales along the axis of the Gulf \(i = 2\) are greater than those across the basin \(i = 1\). For instance, the absolute diffusion is more than two times along than across the GC. This result is expected because the elongated basin geometry limits (allows) the particle motion across (along) the Gulf. The same calculations with the non-rotated velocity components (not shown here) are unable to capture this difference. Note that the standard deviations are rather important, reflecting the inhomogeneity of the buoy data.

Since several drifters have a long life time (up to 90 days), it is possible to increase the number of degrees of freedom by considering subdivisions of buoy tracks whose origin is \(T_i = 10\) days apart. This is a reasonable period after which every subtrack can be considered independent since the Lagrangian timescales are much shorter (Poulain and Niiler, 1989). By proceeding so, the number of segments obtained from the 61 buoys is 244. Using the extended set of segments, the Lagrangian scales are calculated again with the same procedure used with the original set. The values are somewhat smaller, as shown in Table 1 (lower part). The Lagrangian scales along the Gulf are again greater than across the basin, which is consistent with the previous calculation. An important criterion to validate the robustness of the measured Lagrangian scales with segmented trajectories is by varying the period \(T_s\). For small values there will be more segments, but the statistics will be biased if \(T_s\) is of the same order as the Lagrangian timescale. On the contrary, large values of \(T_s\) imply few segments and make the statistics unreliable. For instance, the Lagrangian scales are slightly increased for longer \(T_s\), say, 20 or 30 days; using these times, however, implies only 107 and 57 segments, respectively.

### 3.2. Absolute dispersion

According with Taylor’s theory for homogeneous and stationary turbulence, the dispersion along direction \(i\) is given by

\[
\overline{A_i^2(t)} = 2 \int_0^t dt' \int_0^{t'} R_{ii}(t') dt'. \tag{12}
\]
The asymptotic behavior of (12) for times either smaller or larger than the Lagrangian timescale are, respectively, the quadratic initial dispersion (ballistic regime) and the linear dispersion for longer times (random-walk regime):

$$\bar{A}_i^2(t) \rightarrow R_{ii} t^2 \quad \text{for } t < T_i,$$

$$\bar{A}_i^2(t) \rightarrow 2R_{ii}^2 t \quad \text{for } t > T_i.$$  (13)  (14)

These cases imply a linear increase of diffusivity (the time-derivative of dispersion) for short times, and a constant value for large periods.

Dispersion can be calculated directly from definition (3) but now using the residual buoy positions. Residuals are obtained by extracting the Lagrangian mean to the position of each buoy, as we did for the velocities:

$$\bar{X}_i(t) = x_i(t) - \frac{t}{T} \int_0^T u_i(t') dt'.$$  (15)

By obtaining $\bar{A}_i^2(t)$ in this way, one can compare its time evolution with the predicted limits (13) and (14). The relevance of this comparison is to determine to what extent the measured oceanic dispersion behaves as the case of homogenous, stationary turbulence.

The absolute dispersion (3) is calculated using the 244 segments described in previous subsection. Dispersion curves along and across the Gulf of California are shown in Fig. 4a in a loglog scale. There are two regimes between 0 and 30 days where a power-law can be estimated as

$$\bar{A}_i^2(t) = b_i t^{a_i}.$$  (16)

The important parameter is the exponent $a_i$, which is shown next to the curves. From 0 to 2 days, that is, up to approximately the Lagrangian time scale, both components of dispersion behave in a similar way with a power-law of $a_i(0 < t < 2) \approx 7/4$, somewhat lower than the quadratic, ballistic growth. Between 2 and 30 days the dispersion components diverge. The power-law along the GC is about $a_i(2 < t < 30) \approx 1$, which corresponds with the random-walk regime expected from conventional theory (linear growth). In contrast, a lower value is found across the GC, $a_i(2 < t < 35) \approx 1/2$, so the Taylor theory does not seem to be valid in this case. This result should not be very surprising, however, when taking into account that particle dispersion across the Gulf is saturated due to the limited size of the domain in that direction.

Fig. 4b shows the number of used segments as a function of time. At initial times there is enough statistical information because there is the contribution of the whole set of 244 segments. For subsequent times there are less segments available since buoys with short lifetimes do not contribute anymore. It has a stair-like shape since the interval for generating segments is 10 days. At $t = 30$ days there is still a reasonable number of segments (135), as the smoothness of the dispersion curves suggest. After this time it is considered that there are no sufficient segments to compute the dispersion appropriately.

According with (4) the diffusion is given by

$$K_i(t) = \frac{1}{2} a_i b_i t^{a_i-1},$$  (17)

which in general is time-dependent except for $a_i = 1$ (random-walk diffusion). Thus, a nearly constant diffusivity $K_i$ is only obtained for the direction along the GC between 2 and 30 days. A representative value can be estimated by calculating the time average of (17) during the period of interest:

$$K_i \approx \frac{a_i b_i}{2(T_{fin} - T_{ini})} \int_{T_{ini}}^{T_{fin}} t^{a_i-1} dt = \frac{b_i}{2} \frac{T_{fin}^{a_i} - T_{ini}^{a_i}}{(T_{fin} - T_{ini})},$$  (18)

where $T_{ini} = 2$ days and $T_{fin} = 30$ days. The result is shown in Table 2 (whole year). This value is similar to that obtained with the Lagrangian scales (Table 1).

Seasonal absolute dispersion can be calculated by considering track segments again. The segments correspond to subdivisions of the trajectories that occur entirely within the corresponding season. The seasonal absolute dispersion curves are shown in Fig. 5. Similar to the results using the whole data set, two regimes are identified between 0 and 2 days, and between 2 and 30 days, and the corresponding power-laws are calculated. The results during the first stage are similar in both directions in the four seasons: initial dispersion is isotropic and close to the ballistic regime ($a_i \approx 2$). At longer times, the dispersion components separate each other. The dispersion component across the basin saturates in all seasons ($a_i < 1$). The dispersion along the GC grows almost linear in time during Winter and Spring ($a_i \approx 1$), while being slower in Summer and faster in Autumn. Table 2 presents the estimation of the seasonal diffusivity $K_i$ along the Gulf according with (18).

Some caution should be taken since there are more segments occurring during Spring and Summer, while Autumn might be under sampled (Table 2). Note also that the sum of segments of the four seasons (233) is lower than the sum for the whole year.
(4.2). The reason is because those segments with a first part in one season and another part in the following season are discarded.

4. Two-particle statistics

Relative dispersion (or two-particle statistics) is calculated from the mean-squared separation of particle pairs:

$$D_i^2(t) = \frac{1}{N_0} \sum_{p=q} [x_i^p(t) - x_i^q(t)]^2,$$

(19)

where $N_0$ is the number of pairs and $x_i^0$ the displacements of particles $p$ and $q$ along direction $i$. Relative dispersion measures how a cloud of particles separates in time. Analogously to absolute diffusion, the relative diffusivity is defined as

$$Y_i(t) = \frac{1}{2} \frac{d}{dt} D_i^2.$$

(20)

In Section 4.1 we first discuss different relative dispersion regimes, which depend on the initial scale of separation, and on the forcing injection scale in the flow. Afterwards, we calculate relative dispersion for different initial pair separations and compare with the theoretical dispersion regimes. In Section 4.2, pair dispersion is analyzed by means of the Finite Scale Lyapunov Exponents, which intrinsically take into account the initial separation between particles to estimate the relative diffusivity.

4.1. Relative dispersion

The typical regimes of relative dispersion depend on the initial length scale of pair separations $D_0$, and on the forcing injection scale in the flow $D_i$ (Babiano et al., 1990). These regimes are derived for the relative diffusion as:

$$Y(t, D_0) \propto \epsilon^{1/3} D^2$$

for $D_0 < D_i$

(21)

$$Y(t, D_0) \propto \epsilon^{1/3} \left( \sqrt{D^2} \right)^{4/3}$$

for $D_0 \geq D_i$

(22)

$$Y \propto \text{cte.}$$

for $D_0 \gg D_i$

(23)

In the ocean, $D_i$ can be regarded as the internal Rossby radius of deformation, at which energy is injected to mesoscale eddies by baroclinic instabilities. The first regime belongs to the enstrophy cascade in two-dimensional turbulence, in which enstrophy is transferred at a rate $\chi$ to length scales shorter than $D_i$ (Lin, 1972). It corresponds with an exponential separation of the particles, which are advected by structures larger than the initial separation distance. This is known as non-local dispersion (Babiano et al., 1990; Lumpkin and Elipot, 2010). The second case is the Richardson regime (4/3s law), which occurs when particles are separated a larger distance within the energy cascade range, with $\epsilon$ the small perturbations.

<table>
<thead>
<tr>
<th>Segment</th>
<th>$K_2$ ($10^7$ cm$^2$ s$^{-1}$)</th>
<th># Segments</th>
</tr>
</thead>
<tbody>
<tr>
<td>Whole year</td>
<td>3.42 ± 0.10</td>
<td>244</td>
</tr>
<tr>
<td>Winter</td>
<td>2.43 ± 0.11</td>
<td>48</td>
</tr>
<tr>
<td>Spring</td>
<td>2.37 ± 0.04</td>
<td>62</td>
</tr>
<tr>
<td>Summer</td>
<td>1.50 ± 0.36</td>
<td>85</td>
</tr>
<tr>
<td>Autumn</td>
<td>6.27 ± 1.48</td>
<td>38</td>
</tr>
</tbody>
</table>

Table 2 Absolute diffusivity along the GC ($K_2$) for the full data set and seasonal data, calculated with the time average [18]. Last column presents the number of segments used for each period.

Fig. 5. Seasonal absolute dispersion across (dashed dark lines) and along (dashed gray lines) the GC vs. time. Solid lines and numbers as in Fig. 4. The number of segments used for each season are indicated in Table 2.
rate of energy transfer (Franzese and Cassiani, 2007). The third regime corresponds to particles that are sufficiently far from each other, at a much larger scale than the energy containing eddies, so the diffusivity is approximately constant. In this Fickian regime the particles disperse randomly since there are no flow structures relating one with the other. It must be recalled that these regimes apply for diffusive processes in non-confined flows, that is, without the influence of lateral walls. As for the absolute dispersion results, we shall show that the elongated shape of the GC has a great influence on the measurements of relative dispersion too.

Before presenting the dispersion measurements, let us mention the nature of the buoy pairs. In order to study relative dispersion in the ocean, it is useful to release several drifters initially separated a short distance from each other, and then calculate how the particle pairs separate in time. The data set of the present study was not designed to obtain relative dispersion, so most of drifters were released in different locations at different times. This gives very few “original” pairs (simultaneously released buoys) to carry out a statistically reliable study. A method to increase the number of available pairs consists of considering “chance” pairs formed by buoys that approach each other at some time (LaCasce and Ohlmann, 2003). In the present case, “chance” pairs are a majority. The number of pairs obtained in this way depends on the initial separation and also on time, since the lifetime of the pair elements is not the same in general.

Relative dispersion is directly calculated from (19) for sets of pairs with different initial separations. Let us first try to discern the exponential regime for short initial separations:

$$D_i^2(t) = D_i^2 \exp(\lambda t^2),$$  

with $\lambda$ an $O(1)$ positive constant (Lin, 1972). In general, this regime is difficult to observe with ocean drifters, since there must be enough pairs initially separated a short distance (a few hundreds of meters). This can be achieved only with experiments specifically designed, using initial separations of order 1 km or less (e.g. LaCasce and Ohlmann, 2003; Poje et al. 2014). There are some initial separations of about 5 km in our data set, but most of pairs are initially separated between 10 and 70 km. For the GC the deformation radius $D_o$ is about 35 km, so it seems unlikely to observe the exponential regime. Fig. 6 presents the measurements of relative dispersion in a logarithmic scale in order to visualize a probable linear relationship of the form

$$\ln D_i^2 = \ln D_i^2 + ct,$$

with $c = \lambda t^2$ the growth rate. The dispersion across and along the GC are presented separately. Each curve is calculated with a special class of particle pairs initially separated a distance $D_i$ within a range of 10 km (10–20 km, 20–30 km, and so on). The short straight lines over the dispersion curves are the best fit to (25) during the first 2.5 days, and the number indicates the slope $c_i$. Note that the growth rates are very similar in both directions regardless of the initial separation classes, and that the growth rates diminish for larger initial separations.

Although these plots seem to represent an exponential regime, the results should be taken with some caution. First, it must be noted that the period at which the exponential curves are fitted (2.5 days) is relatively short (only contains 11 data points). Secondly, the number of pairs might be insufficient in some cases, specially for the first class (10–20 km). The decreasing number of pairs as time progresses is shown in Fig. 7, where it is verified that the shortest initial separation class (panel a) only contains 17 pairs at initial times, which is a very reduced number.

We now evaluate the dispersion behavior when particles separate a longer distance. Fig. 8 shows the relative dispersion again, but now in log–log plots with different initial separations and for longer times. The dispersion across the Gulf is saturated at about day 3 or 4, while dispersion along the GC continues growing during longer times. Saturation is obviously associated with the short dimensions across the GC, so the dispersion in this direction has an erratic behavior after day 4. In contrast, dispersion along the GC grows in time during 20–30 days more. For initial separations shorter than the deformation radius, $D_o < D_i$ the dispersion would enter into the Richardson regime, whose asymptotic behavior is the well-known $r^2$ law (Babiano et al., 1990)

$$D_i^2(t) \propto t^2.$$  

Despite $D_o$ in our data set is of the same order or larger than $D_i$, we can look for the presence of the Richardson regime along the GC by fitting a power-law of the form

$$D_i^2 = \epsilon t^2,$$

(with $\epsilon$ a constant with appropriate units) and examining the value of $d$. The best fit is also presented in Fig. 8. In all cases it is found that the exponent is around 3/2, so the growth of relative dispersion is slower than the $r^2$ asymptotic law. Similar results have been found in previous studies with oceanographic data: LaCasce and
Ohlmann (2013) found $d/C_2^2$ in the Gulf of Mexico, while Haza et al. (2008) reported $d/C_1^9$ in the Adriatic Sea. Note that the power of time increases for shorter initial separations. This is consistent with Babiano et al. (1990), as mentioned above. Since the shortest $D_0$ is $30–40$ km, for which the exponent is $1.74$, then it is not expected to recover the $t^3$ law.

Does this result proves the non-existence of a Richardson regime? In order to reevaluate this, we consider the $4/3$s Richardson law between the diffusivity and the relative dispersion:

$$Y_i(t, D_0) = \frac{1}{2} \frac{d}{dt} \sqrt{D_i} = \beta \epsilon^{1/3} \left( \sqrt{D_i} \right)^{4/3},$$  \hspace{1cm} (28)

with $\beta$ an $O(1)$ non-dimensional constant. Then we follow the procedure of Ollitrault et al. (2005), who integrated (28) and obtained the following form of the relative dispersion (in our case, for the direction along the GC):

$$\sqrt{D_i} = (b + at)^3,$$  \hspace{1cm} (29)

with $b = (D_0/2)^{1/3}$ and $a = \frac{1}{2} \beta \epsilon^{1/3}$. This expression indicates that the dispersion grows as a third-order polynomial, and not just $t^1$. To explore whether our results fit well with this expression, we take the cubic root of (29),

$$\sqrt{D_i}^{1/3} = b + at,$$  \hspace{1cm} (30)

and then $a$ and $b$ are calculated by fitting a straight line. This process is repeated for each class of initial separations; the results are shown in Table 3. The $a$-values allow an estimation of the (cubic root of the) rate of energy transfer $\epsilon$ (last column), which has a similar order as that reported in previous studies (see Ollitrault et al., 2005 and references therein). The $b$-values are close, though somewhat smaller, than the corresponding $(D_0/2)^{1/3}$.

Then, the time coordinate is shifted as $^\wedge t = t + b/a$ (Ollitrault et al., 2005) which yields

$$\sqrt{D_i} = a^{3} t^{3}.$$ \hspace{1cm} (31)

Fig. 9a shows the dispersion curves for four initial separation classes as a function of the shifted-time $^\wedge t$. The curves indeed grow as $^\wedge t^1$, corroborating expression (31). The data of the different initial separations collapse into one single curve by dividing over the corresponding $a^3$-value. This is presented in Fig. 9b, where the average slope (in the log–log scale) is indicated.

A similar situation can be calculated for the diffusivity by taking (half) the time-derivative of (29), and shifting the time as before:

$$Y_2 = \frac{3}{2} a^2 t^2.$$ \hspace{1cm} (32)

Fig. 9c presents the corresponding curves of diffusivity in terms of separations $Y_2$ vs. $(D_i/2)^{1/2}$. $Y_2$ is calculated as a function of time by using (32), while $(D_i/2)^{1/2}$ is directly measured from the buoy data. The results confirm that diffusion grows as stated by expression (28). The curves collapse when plotting $Y_2/(3a/2)$ (Fig. 9d), obtaining a curve whose average slope (in the log–log scale) is pretty close to $4/3$. As pointed out by Ollitrault et al. (2005), the plots in Fig. 9 do not prove that the Richardson regime is obeyed.
by the pairs, but that the measured relative dispersion reasonably agrees with this behavior. These results are further discussed below.

4.2. Finite scale Lyapunov exponents

Relative dispersion measures the average particle separations that are different initially, which may obscure its time dependence (LaCasce, 2008). A different method to estimate the dispersion of a set of particles is given by the Finite Scale Lyapunov Exponents (FSLE) introduced by Artale et al. (1997) and Aurell et al. (1997). The FSLE are commonly used for measuring dispersion properties (Lacorata et al., 2001) or to identify Lagrangian structures (Boffeta et al., 2001; d’Ovidio et al., 2004). A remarkable advantage of this approach is that the diffusivity is naturally identified as a function of the particle separations, as proposed by Richardson (1926). The main idea is to average times of separation instead of averaging distances.

The basic procedure behind the FSLE is to choose a particle pair separated by a distance \( \delta \), and then calculate the time \( \tau(\delta) \) at which the particles separate a distance \( r\delta \), where \( r \) is an \( O(1) \) positive real number such that \( r > 1 \). Here we use \( r = \sqrt{2} \) (Lacorata et al., 2001). In general \( \tau(\delta) \) depends on the initial separation. The calculation is repeated for all available pairs. The FSLE are defined as

\[
\lambda(\delta) = \frac{1}{\langle \tau(\delta) \rangle} \log r,
\]

where \( \langle \tau(\delta) \rangle \) is the average of all the separation times for a given \( \delta \). Evidently, \( \lambda(\delta) \) is the rate of divergence between two particle positions, as the conventional Lyapunov exponent, but now calculated with finite separations (d’Ovidio et al., 2004).

The aim, then, is to find the relation \( \lambda(\delta) \) vs. \( \delta \). The theoretical regimes are described by Artale et al. (1997) and Lacorata et al. (2001). In terms of the size of the energy containing eddies \( L_t \):

\[
\lambda(\delta) \approx \lambda_0 \quad \text{for } \delta \ll L_t, \quad \text{with } \lambda_0 \text{ constant}
\]

\[
\lambda(\delta) \propto \delta^\gamma \quad \text{for } L_t \ll \delta, \quad \text{with } -2 < \gamma < 0
\]

\[
\lambda(\delta) \propto \delta^{-2/3} \quad \text{for } L_t \ll \delta.
\]

For very short scales in the first regime the FSLE are constant, which implies an exponential pair separation, that is, Lagrangian chaos. For large scales in the third case, particle velocities are uncorrelated and the standard diffusion process is recovered. The intermediate regime corresponds to the size of the energy containing eddies, which might be present in oceanic observations. The Richardson regime corresponds with exponent \( \gamma = -2/3 \).

The calculation of the FSLE is performed over the available pairs constructed from the 61 buoys of the original data set. The particles belonging to one pair are simultaneous (by definition), but

---

**Table 3**

Fitting parameters \( a \) and \( b \) in the linear expression (30) for each separation class. Last column presents \( \beta \sqrt{a/3} \) in appropriate units.

<table>
<thead>
<tr>
<th>Separation class (km)</th>
<th>( a ) (km(^2) d(^{-1}))</th>
<th>( b ) (km(^3)d(^{-1}))</th>
<th>( \beta \sqrt{a/3} ) (m(^{3/2}) s(^{-1}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>30–40</td>
<td>1.76</td>
<td>8.66</td>
<td>3.0 \times 10^{-3}</td>
</tr>
<tr>
<td>40–50</td>
<td>1.40</td>
<td>10.75</td>
<td>2.4 \times 10^{-3}</td>
</tr>
<tr>
<td>50–60</td>
<td>1.33</td>
<td>12.33</td>
<td>2.3 \times 10^{-3}</td>
</tr>
<tr>
<td>60–70</td>
<td>0.98</td>
<td>15.53</td>
<td>1.7 \times 10^{-3}</td>
</tr>
</tbody>
</table>

---

*Fig. 8.* Relative dispersion vs. time in log–log plots for different initial separations (panels a to d): across the GC (dotted dark line), and (b) along the Gulf (dotted gray line). The numbers indicate the slope \( d \) when fitting expression (27) to the dispersion curves along the GC (black lines). Dashed lines indicate the 90% confidence limits.
pairs might be simultaneous or not. This was a limitation for the calculation of relative dispersion, but it is not for the FSLE, with which the number of degrees of freedom is sensibly increased. The initial separations are chosen within \( n r_0 \) and width \( r_0 \), so the first bin includes all particle pairs with an initial separation between \( r_0 - r_0/2 \) and \( r_0 + r_0/2 \), the second bin between \( 2r_0 - r_0/2 \) and \( 2r_0 + r_0/2 \), and so on, up to the interval between \( (n - 1)/2 r_0 \) and \( (n + 1)/2 r_0 \).

Fig. 10a and b shows the FSLE \( \lambda(\delta) \) for two ways of binning, \( r_0 = 10 \) and 15 km. The chaotic regime (constant \( \lambda \) for small \( \delta \)) is not captured since the shortest initial separations are 5 km and 7.5 km, respectively. This regime is difficult to observe with the present data set, since the buoy deployment was not designed to have small initial separations. Pairs with separations shorter than 5 km are available, but they are too few to be statistically reliable. The standard diffusion regime (large \( \delta \)) is not obtained either, perhaps due to the finite size of the Gulf or to the lack of data at large scales. However, an intermediate diffusion regime with \( \gamma \sim -2/3 \) can be clearly identified by fitting a straight line to the log-log plots. This corresponds to the Richardson regime in the GC.

In order to analyze this result, it is important to verify the number of available pairs for each bin. This is presented in Fig. 11a for bins with \( r_0 = 5, 10, \) and 15 km. The Richardson regime is measured when considering data up to \( \delta \sim 140 \) km approximately, which is the distance at which the number of available pairs starts to decrease. Note that this distance corresponds with the average width of the GC. Now we can estimate a diffusivity coefficient for different length scales as \( Y(\delta) = 2\lambda \delta^2 \) (Artale et al., 1997). The corresponding curves for the three different binnings is shown in Fig. 11b. When fitting a straight line with the \( r_0 = 10 \) km bins, we obtain a very precise representation of the 4/3s Richardson law.

5. Discussions

Absolute dispersion obeys a nearly ballistic regime during 2 days, approximately, in the directions across and along the GC. After that time, it is saturated across the Gulf due to the presence of the eastern and western boundaries. The exact ballistic behavior is quadratic; in our measurements, however, the time exponent seems to be somewhat smaller than 2 in both directions, either for the full set of drifters (~7/4) or at a seasonal scale (except on Autumn). There are no clear reasons for this discrepancy. Besides associated errors in the measurements, one possibility is the presence of anomalous dispersion driven by the dominant influence of hyperbolic or elliptical regions in the context of two-dimensional turbulent flows (Elhmaïdi et al., 1993). In a similar vein, Boffetta et al. (1997) have shown that anomalous exponents arise as a consequence of averaging groups (“populations”) of walkers with different decorrelation times. Such differences can be attributed to the presence of coherent structures, and can lead to power-laws \( t^x \) with \( 1 < x < 2 \).

The dispersion along the GC continues growing linearly in time, \( D_s^t \propto t^\alpha \) with \( \alpha \sim 1 \), during 25–30 days more. This behavior corresponds to the random walk regime, expected for long times, with a constant diffusion coefficient \( K_s \sim 3.5 \times 10^7 \) cm² s⁻¹. Similar results are measured for Winter and Spring seasons, while
dispersion in Summer seems to be sub-diffusive ($a < 1$), and Autumn super-diffusive ($a > 1$). We consider that diffusivity in the GC is influenced by the presence of mesoscale vortices that occupy the width of the basin during part of the year. Some of these features are generated during the Summer at the southern and central regions, and persist during several months (Figueroa et al., 2003; Lavín et al., 2013, 2014). At the northern part there is a semi-permanent seasonal gyre (Lavín et al., 1997). When traveling along the Gulf (in either direction), the buoy dispersion might be inhibited by these structures when being captured during some days or weeks. In some cases the trapping fully prevents the dispersion of the particles or even reverse their direction of motion. This might explain the sub-diffusive behavior during the Summer. Across the Gulf, in contrast, mesoscale vortices might be an efficient mechanism to transport particles between the continent and Baja California, thus promoting the saturation of dispersion in this direction in a few days.

The specific role of mesoscale vortices along the GC should be elucidated in a more comprehensive study. One possible methodology is that used by Veneziani et al. (2004) to analyze Lagrangian data in the Northwest Atlantic Ocean. Essentially, these authors decompose the mesoscale turbulent field in looping and nonlooping motions, the former being characterized by a measurable spin obtained by means of a Lagrangian stochastic model. In this way, regions with dominant cyclonic or anticyclonic loopers can be identified, and then determine their influence on the dispersion measurements.

Diffusivity along the GC is comparable with measurements in other regions. The values reported by Poulain and Niiler (1989) for the California Current System (using 595 segments of their buoy trajectories) are $3.4$ and $4.3 \times 10^7$ cm$^2$ s$^{-1}$ along the zonal and meridional directions, respectively. Colin de Verdiere (1983) found zonal and meridional diffusivities of $2.3$ and $1.7 \times 10^7$ cm$^2$ s$^{-1}$ at the eastern North Atlantic. In oceanic systems dominated by the presence of intense currents, diffusion can be as large as $10^8$ cm$^2$ s$^{-1}$, as measured in the Algerian current (Salas et al., 2001). It is illustrative to compare the dispersion results in the GC with previous findings in the Adriatic Sea, which has a similar elongated shape though being slightly shorter ($\sim 800$ km) and wider ($\sim 200$ km). The measured diffusivities by Falco et al. (2000) are $1$ and $3 \times 10^7$ cm$^2$ s$^{-1}$ for the zonal and meridional directions, respectively. One of the major differences between both semi-enclosed seas is that the circulation in the Adriatic is mostly cyclonic, dominated by two basin-wide gyres located in the middle

![Fig. 10.](image-url) Finite Scale Lyapunov Exponents (FSLE) vs. separations $\delta$ for two ways of binning: (a) $r_0 = 10$ km and (b) $r_0 = 15$ km. The solid line is the best fit considering the first 14 (140 km) and 9 (145 km) data points, respectively. The dashed lines are the associated errors, and the numbers indicate the power-law exponent $-\gamma$.

![Fig. 11.](image-url) (a) Number of pairs as a function of distance $\delta$ for three different ways of binning: $r_0 = 5$ km (circles), $r_0 = 10$ km (stars), and $r_0 = 15$ km (plus signs). The vertical dashed line indicates $\delta = 140$ km. (b) Diffusivity for the same three ways of binning. The solid line is the best fit for the $r_0 = 10$ km data, considering the first 14 (140 km), and the number indicates the corresponding power law. Note that $Y \approx \delta^{4.3}$, which corresponds with the Richardson regime.
and southern regions (Lacorata et al., 2001), and with intense coastal currents along the long sides that exert a strong influence in the dispersion properties (Haza et al., 2008). Furthermore, a train of mesoscales vortices is not as clear as the one observed in the GC. However, these are just preliminary observations. We suggest that, given the similarity between the two elongated seas, a comparative study would be useful to gain further insight on the dispersion problem in an elongated, semi-enclosed domain.

Relative dispersion was measured, despite the data set used in the present study was not designed for this purpose. We first looked for the exponential growth of relative dispersion for particles that are initially separated a short distance (compared with the deformation radius of about $D_t \sim 35$ km), according with the enstrophy cascade regime in two-dimensional turbulence. However, the results are not conclusive because there are very few pairs with initial separations shorter than $D_t$. There is an open discussion regarding the validity of the non-local dispersion theory in oceanic flows. Although some authors have reported the exponential growth for initial separations as short as 1 km (LaCasce and Ohlmann, 2003; Koszalka et al., 2005) or even 7.5 km (Ollitrault et al., 2005), recent studies specifically designed to resolve submesoscale motions have not registered this regime (Lumpkin and Elipot, 2010; Poje et al., 2014). Besides methodological arguments related with the use of chance pairs, Lumpkin and Elipot (2010) discussed that the reason might be associated with the injection of energy into the submesoscale range associated with frontal and inertial instabilities, and more in general to the surface quasi-geostrophic dynamics. We cannot elaborate more in this study because with the present data set we cannot accurately resolve submesoscale flows.

Secondly, we analyzed the relative dispersion for initial separations of the same order or larger than $D_t$. Across the GC, relative dispersion is saturated again. The saturation of dispersion across the Adriatic was also reported by Lacorata et al. (2001). Along the Gulf, we explored the possible presence of the classical asymptotic regime for homogeneous turbulence, which says that dispersion increases as $t^3$ (Richardson regime). However, it is found that relative dispersion grows as $t^{1.5}$, approximately. It is argued again that this slower behavior might be due to the presence of persistent mesoscale vortices that occupy the width of the Gulf, as previously described. Some other reports have also registered dispersion growths slower than the $t^3$ law. For instance, Okubo (1971) showed that the dispersion measured in several studies is proportional to $t^2$ with $2 < x < 3$ (most of such experiments were carried out in systems with smaller length scales). Using a high resolution model of the Adriatic Sea circulation, Haza et al. (2008) calculated synthetic trajectories from which relative dispersion was shown to evolve as $t^{1.9}$, approximately. LaCasce and Ohlmann (2003) found a growth of $t^{2.2}$ in the Gulf of Mexico. A possible explanation for the break down of the cubic law was provided by Babiano et al. (1990), who found that asymptotic regimes in homogeneous turbulence, such as the $t^3$ law, can only be reached when the initial separation is small, which is not the case in the present study. However, the experiments of LaCasce and Ohlmann (2003) were designed with initial separations of the order of 1 km and yet did not measured a $t^3$ growth. As pointed out by Franzese and Cassiani (2007), the evidence for this regime is quite controversial in both oceanic and atmospheric measurements.

The Richardson regime can be examined by integrating the $4/3$s Richardson law and considering a time shift, as in Eq. (31). This procedure was apparently first noted by Ollitrault et al. (2005).

Essentially, this expression indicates that the dispersion grows as a general, third-order polynomial, and not just $t^3$. Franzese and Cassiani (2007) discussed that the $t^3$ scaling breaks down in the vicinity of a finite source of particles, because the statistical properties of the diffusion are affected by the different initial conditions (particle separations) associated with the finite size of the source. In order to take into account this effect, Franzese and Cassiani (2007) derived a similar time integration of the Richardson law and obtained an expression equivalent to (31). In terms of our own notation, they denoted the time shift $t_s = b/a = (3/2)(\langle D_t^2 \rangle / \epsilon) ^ {1/3}$, which is interpreted as the time required for a set of particles emerging from a point source to expand to the mean square size $\langle D_t^2 \rangle$. Here we showed that relative dispersion along the GC indeed grows as the cube of the shifted time, $\langle \tilde{D}_t^2 \rangle \propto t^3$.

The analysis with the FSLE is consistent with the findings of the Richardson regime. This behavior seems to be very robust, since it is obtained by using different ways of binning the separation distances between particles. This result supports the notion that the use of FSLE is a more effective method to determine relative dispersion than the direct use of formula (19), which has to be calculated for different sets of pairs with different initial separations. In contrast, the dependence of the diffusion coefficient with respect to the particle separations arises naturally from the calculation of the FSLE.

6. Conclusions

In this paper we examined the Lagrangian transport and dispersion of surface drifters along the Gulf of California. Analysis of one and two-particle statistics has been carried out by using the trajectories of 61 buoys released during 2004–2006. For the former case, we calculated direct measurements of absolute dispersion and Lagrangian scales based on the classical theory for homogeneous turbulence. For the two-particle statistics, we measured relative dispersion of pairs of particles, as well as Finite Scale Lyapunov Exponents.

The dispersion properties have been discussed for the transversal and longitudinal directions of the GC. For short times (a few days), absolute dispersion is nearly ballistic in both directions, while relative dispersion suggested exponential separations of particle pairs (but no conclusive results were found for this regime). At longer times (20–30 days), in contrast, the elongated character of the basin implies that dispersion along the Gulf strongly differs from that across the GC. In the transversal direction, both absolute and relative dispersion are limited by the long sides of the basin (the continental Mexican coast at the eastern side and the Peninsula of Baja California to the west). Thus, individual particles and pairs are preferentially dispersed along the GC: absolute dispersion obeys a random walk behavior, while relative dispersion follows a Richardson regime.

A second essential result is that the typical circulation of the region, characterized by the emergence of mesoscale vortices during some periods of the year, apparently plays a fundamental role in the dispersion. For instance, the northern part is relatively shallow and the surface circulation is dominated by the seasonal gyre, while the southern region is much deeper and the circulation is strongly affected by mesoscale eddies and the influence of the Pacific Ocean. When particles travel along the Gulf, they might become trapped within these structures, preventing or delaying further displacements. The vortices are also an important mechanism to translate particles across the Gulf, between the Peninsula and the continent.

One further consideration is that the present study provides a global view of the dispersion in the GC. More detailed analyses should be carried out in specific regions and seasons of the year, because the characteristics of the Gulf circulation have great variations along its elongated shape, and change sensibly during the
Another suggestion is to perform comparative studies with other basins with a similar morphology, such as the Adriatic Sea. A useful design for future experiments should contemplate the deploy of simultaneous drifters with very short initial separations (e.g. 1 km) and long enough tracking times (months), in order to investigate the presence of the exponential growth of relative dispersion. The positive verification of this regime would imply that the dispersion can be explained by means of the enstrophy cascade, whereas the negative result suggests that submesoscale motions are rather important (Lumpkin and Elipot, 2010).

Acknowledgments

The author gratefully acknowledges the support of Dr. Miguel Fernando Lavín Peregrina (1951–2014), who kindly provided the buoy data during the first stages of this research. Miguel was an excellent teacher, a thorough researcher and a very generous person, who devoted much of his life to unravel the physical oceanography of the Gulf of California.

References


