Lagrangian Drifter Dispersion in the Surfzone: Directionally-Spread Normally- Incident Waves

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April 17, 2008

ABSTRACT

Lagrangian drifter statistics in a surfzone wave and circulation model are examined and compared to single- and two-particle dispersions statistics observed on an alongshore uniform natural beach with small normally-incident, directionally spread waves. Drifter trajectories are modeled with a time-dependent Boussinesq wave model that resolves individual waves and parameterizes wave breaking. The model reproduces the cross-shore variation in wave statistics observed at three cross-shore locations. In addition, observed and modeled Eulerian binned (means and standard deviations) drifter velocities agree. Modeled surfzone Lagrangian statistics are similar to those observed. The single-particle (absolute) dispersion statistics are well predicted, including non-dimensionalized displacement pdfs and the growth of displacement variance with time. The modeled relative dispersion and scale-dependent diffusivity is consistent with the observed and indicates the presence of a 2D turbulent flow field. The model dispersion is due to the rotational components of the modeled velocity field, indicating the importance of vorticity in driving surfzone dispersion. Modeled irrotational velocities have little dispersive capacity. Surfzone vorticity is generated by finite crest-length wave-breaking that, on the alongshore uniform bathymetry, results from a directionally spread wave field. The generated vorticity then cascades to other length-scales as in 2D turbulence. Increasing the wave directional spread results in increased surfzone vorticity variability and surfzone dispersion. Eulerian and Lagrangian analysis of the flow indicate that the surfzone is 2D turbulent-like with an enstrophy cascade for length-scales between approximately 5–10 m and an inverse-energy cascade for scales 20–100 m. The vorticity injection length-scale (the transition between enstrophy to inverse-energy cascade) is a function of the wave directional spread.

1. Introduction

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Terrestrial runoff dominates urban pollutant loading rates (Schiff et al. 2000). Often draining directly onto the shoreline, runoff pollution degrades surfzone water quality, leading to beach closures (e.g., Boehm et al. 2002). Runoff increases the health risks (e.g., diarrhea and upper respiratory illness) to ocean bathers (Haile et al. 1999), and contains both human viruses (Jiang and Chu 2004) and elevated levels of fecal indicator bacteria (Reeves et al. 2004). Surfzone mixing processes disperse and dilute such pollution. The surfzone and nearshore are vital habitat to ecologically and economically important species of marine fish (e.g., Romer and MacLachlan 1986) and invertebrates (e.g., Lewin 1979). The same surfzone dispersal processes likely affect nutrient availability, primary productivity, and larval dispersal (e.g., Talbot and Bate 1987; Denny and Shibata 1989). Understanding surfzone Lagrangian dispersion processes is important to predicting the fate (transport, dispersal, and dilution) of surfzone tracers whether pollution, bacteria, larvae, or nutrients.

Previous surfzone dispersion studies have generally tracked fluorescent dye (Harris et al. 1963; Inman et al. 1971; Grant et al. 2005; Clarke et al. 2007) resulting in estimated “eddy” diffusivity magnitudes that vary considerably. These studies have difficulties in detailed dye tracking and are based on single realizations. Surfzone Lagrangian drifters also are used to study dispersion. Johnson and Pattiaratchi (2004) used the spreading rate of multiple drifters to estimate scale-dependent relative diffusivities in the surfzone on a beach with a dominant rip-current circulation feature. For approximately 10-50 m separations, relative diffusivities between $1.3 \times 10^{-2}$ to $3.9 \times 10^{-2}$ m$^2$s$^{-1}$ were reported.

Two days of surfzone drifter dispersion observations on an alongshore uniform beach were reported by Spydell et al. (2007). The first day had small normally incident waves with weak mean currents and the second day had large obliquely incident waves driving a strong alongshore current. Absolute and relative Lagrangian statistics were presented for both days. On the first day the observed drifter dispersion had properties similar to a two-dimensional (2D) turbulent fluid and the scale-dependent relative diffusivity suggested the presence of a surfzone eddy (vorticity) field with a range of length-scales spanning 5-50 m (Spydell et al. 2007). The lack of any mean currents precludes sheared currents as the source of this eddy field and finite crest length breaking waves was hypothesized to generate the vertical vorticity (Pergerine 1998). This vorticity could then cascade to other length-scales analogous to the vorticity dynamics of 2D turbulence. On an alongshore uniform beach, finite breaking crest length is the result of non-zero wave directional spread $\sigma_\theta$ (e.g., Kuik et al. 1988), that is incoming waves with a variety of angles.

Accurately modeling and diagnosing surfzone dispersion requires resolving dynamics on a wide range of time-scales from surface gravity waves (few seconds) to very low frequency vortical motions (1000 s) and length-scales from a few meters to many multiple surfzone widths (1000 m). In addition, representing the effects of finite crest length wave-breaking is hypothesized to be important. Time-dependent Boussinesq wave models (e.g., Nwogu 1993; Wei et al. 1995), that simulate wave breaking with an eddy viscosity term in the momentum equations (e.g., Chen et al. 1999; Kennedy et al. 2000) associated with the front-face of steep (breaking) waves, fit these requirements. These types of Boussinesq models reproduce observed wave height variation across the surfzone in the laboratory (Kennedy et al. 2000) and field (Chen et al. 2003). In addition to representing the 2D (horizontal) nature of shoaling and breaking waves (Chen et al. 2000), Boussinesq model simulations with directionally spread waves give rise to a rich surfzone eddy field with vorticity variability over a range of scales (Chen et al. 2003; Johnson and Pattiaratchi 2006).

Here, the question of whether surfzone vorticity and the resulting Lagrangian dispersion is consistent with a “2D turbulent” fluid is examined with Eulerian and Lagrangian statistics. In forced 2D turbulence,
energy is injected at a particular length-scale. The resulting turbulent eddies then cascade to other length-scales following 2D vorticity dynamics resulting in two classifiable regimes, the inverse-energy and enstrophy cascade regions. In the inverse-energy cascade of 2D (and also in inertial subrange of 3D) turbulence (i.e., spatial scales larger than the turbulent injection scale), the Eulerian signature is an \( E \sim k^{-5/3} \) velocity wavenumber spectrum whereas the Lagrangian signatures relate to particle separation statistics. Specifically, the variance of particle separations depends on the cube of time \( D^2 \sim t^3 \), the probability density function (pdf) of separations is non-Gaussian \( P(r) \sim \exp(-|r|^{2/3}) \), and the relative diffusivity depends on the separation \( \kappa \sim r^{4/3} \). These scalings are collectively considered Richardson’s laws which were first obtained empirically for atmospheric data (Richardson 1926). The theoretical basis of these scalings (excluding the pdf shape) derive from dimensional arguments (Obukhov 1941a,b; Batchelor 1950). These laws have subsequently been observed in DNS simulations of 2D (Boffetta and Sokolov 2002b) and 3D (Boffetta and Sokolov 2002a) turbulence and in laboratory experiments of 2D turbulence (Jullien et al. 1999). Furthermore some oceanic observation are consistent with Richardson’s laws (e.g., Stommel 1949; Okubo 1971).

In the enstrophy cascade of 2D turbulence (i.e., spatial scales smaller than the turbulent injection scale), the energy spectrum is given by \( E(k) \sim k^{-3} \) with separation variance growing exponentially \( D^2 \sim \exp(t) \) and the relative diffusivity strongly scale-dependent \( \kappa \sim D^2 \). These Lagrangian enstrophy cascade laws were originally motivated by atmospheric data (Lin 1972) and later observed in laboratory experiments of 2D turbulence (Jullien 2003).

Here the Spydell et al. (2007) day one surfzone drifter observations (Section 2) are simulated with a Boussinesq model, described in Section 3. Lagrangian single- and two-particle statistics are described in Section 4. Model-data comparison of the Eulerian wave and current statistics give good agreement (Section 5), indicating that the surfzone processes are reasonably represented by the model. Lagrangian absolute- and relative dispersion model-data comparison is reported in Section 6. Both absolute and relative dispersion statistics compare well although the magnitude of the observed relative dispersion is larger and scales more slowly with time than the modeled relative dispersion. Both enstrophy and inverse-energy cascades are inferred from the modeled Lagrangian statistics.

The Boussinesq model is used to diagnose the underlying processes leading to the modeled dispersion. Model velocity fields are decomposed into irrotational and rotation components and drifters are advected within each velocity field (Section 7). At times \( > 30 \) s, (absolute and relative) dispersion is dominated by rotational velocities, indicating the importance of vorticity even on an alongshore uniform bathymetry with weak alongshore currents. The vorticity generation mechanism is the non-zero curl of the force imparted by the Boussinesq model wave breaking formulation. This mechanism is identical to the alongshore gradients in breaking wave dissipation discussed in Peregrine (1998) and requires a directionally spread wave field to create finite breaking crest lengths. Boussinesq model simulations with varying incoming wave directional spread \( \sigma_{\theta_0} \) are used to investigate the relationship between \( \sigma_{\theta_0} \), the fluctuating vorticity field, and the resulting surfzone drifter dispersion (Section 8). Eulerian analysis of the model data at various \( \sigma_{\theta_0} \) reveals regimes of both enstrophy and inverse-energy cascades, with the length-scale separating the two regimes depending upon \( \sigma_{\theta_0} \), i.e., the length-scale of vorticity injection is \( \sigma_{\theta_0} \) dependent. This vorticity then freely evolves and cascades to other length-scales in a 2D turbulence-like fashion. The results are summarized in Section 9.
2. The Observations

Observations of surfzone drifter dispersion were acquired on Nov 3, 2004 at Torrey Pines beach in San Diego, CA with small normally incident directionally-spread waves and weak mean currents. These observations are reported upon in detail in Spydell et al. (2007) and are briefly described here. The cross- and alongshore coordinates are $x$ and $y$, with $x = 0$ m near the mean shoreline and $x$ increasing negatively offshore. Locally, the bathymetry was nearly alongshore uniform. The bathymetry alongshore uniformity statistic (Ruessink et al. 2001) $\chi^2 = 0.0036$ in the inner-surfzone region, an order of magnitude smaller than that found to create alongshore non-uniform circulation (Ruessink et al. 2001; Feddersen and Guza 2003). Three Sontek Triton Acoustic Doppler Velocimeters (ADV), sampling at 2 Hz, were deployed on a cross-shore transect with sensing volumes 0.8 m above the bed and were used to estimate wave statistics such as significant wave height $H_s$, mean wave angle $\bar{\theta}$, and wave directional spread $\sigma_\theta$ (Appendix A).

Drifter deployments were conducted over 6 hours with roughly stationary wave, current, and tide conditions (Spydell et al. 2007). There were 9 separate releases of 9 drifters on a cross-shore transect. A total of 77, approximately 1000 s long, drifter trajectories passed quality control. The freely floating, impact resistant, GPS-tracked surfzone drifters are 0.5-m tall cylinders with most of their volume below the water line. A horizontal disc at the bottom of the body tube dampens vertical motions in the waves, allowing broken waves to pass over the drifter without pushing or “surfing” it ashore. Drifter GPS positions are internally recorded at 1 Hz with absolute position error of about ±4 m (George and Largier 1996). Post-processing using carrier phase information reduces the absolute error to ±1 m (Doutt et al. 1998). Technical description of the drifters and its response are found in Schmidt et al. (2003).

3. The Boussinesq Wave and Current Model

A time-dependent Boussinesq wave model similar to FUNWAVE (e.g., Chen et al. 1999), which resolves individual waves and parameterizes wave breaking is used to numerically simulate velocities and sea surface height in the surfzone. The Boussinesq model equations are similar to the nonlinear shallow water equations but include higher order dispersive terms (and in some derivations higher order nonlinear terms). Here the equations of Nwogu (1993) are implemented. The equation for mass (or volume) conservation is

$$\frac{\partial \eta}{\partial t} + \nabla \cdot [(h + \eta)u] + \nabla \cdot \mathbf{M}_d = 0,$$

(1)
where \( \eta \) is the instantaneous free surface elevation, \( t \) is time, \( h \) is the still water depth, \( u \) is the instantaneous horizontal velocity at the reference depth \( z_r = -0.531h \), where \( z = 0 \) at the still water surface. The dispersive term \( M_d \) in (1) is
\[
M_d = \left( \frac{z_r^2}{2} - \frac{h^2}{6} \right) h \nabla (\nabla \cdot u) + (z_r + h/2) h \nabla [\nabla \cdot (hu)].
\] (2)

The momentum equation is
\[
\frac{\partial u}{\partial t} + u \cdot \nabla u = -g \nabla \eta + F_d + F_{br} - \frac{\tau_b}{(\eta + h)} - \nu_{bi} \nabla^4 u,
\] (3)
where \( g \) is gravity, \( F_d \) are the higher order dispersive terms, \( F_{br} \) is the breaking terms, \( \tau_b \) is the instantaneous bottom stress and \( \nu_{bi} \) is the hyperviscosity for the biharmonic friction (\( \nabla^4 u \)) term. The dispersive terms are (Nwogu 1993)
\[
F_d = - \left[ \frac{z_r^2}{2} \nabla (\nabla \cdot u_t) + z_r \nabla (\nabla \cdot (hu_t)) \right],
\]
and the bottom stress is parameterized with a quadratic drag law
\[
\tau_b = c_d |u| u,
\]
with the non-dimensional drag coefficient \( c_d \).

Following Kennedy et al. (2000), the effect of wave breaking on the momentum equations is parameterized as a Newtonian damping where
\[
F_{br} = (h + \eta)^{-1} \nabla \cdot \left[ \nu_{br}(h + \eta) \nabla u \right],
\]
where \( \nu_{br} \) is the eddy viscosity associated with the breaking waves.\(^1\) The breaking eddy viscosity is given by
\[
\nu_{br} = B \delta^2 (h + \eta) \frac{\partial \eta}{\partial t},
\] (4)
where \( \delta \) is a constant and \( B \) is a function of \( \eta_t \) and varies between 0 and 1. When \( \eta_t \) is sufficiently large (i.e., the front face of a steep breaking wave) \( B \) becomes non-zero. The Kennedy et al. (2000) expression for \( B \) is used here. The wave breaking parameter choices are similar to the ones used by Kennedy et al. (2000) to model laboratory breaking waves and Chen et al. (2003) for modeling laboratory and field wave heights and alongshore currents. The model results are not overly sensitive to these choices.

The model equations (e.g., (1) & (3)) are 2nd order spatially discretized on a C-grid (Harlow and Welch 1965) and time-integrated with a third-order Adams-Bashforth (Durran 1991) scheme. The model extent is 482 m in the cross-shore, excluding sponge layers (see Fig. 1) and 2000 m in the alongshore. The alongshore boundary conditions are periodic. The cross-shore and alongshore grid spacing is 1 m and 2 m, respectively. The model time step is \( \Delta t = 0.01 \) s. The bathymetry is alongshore uniform and equal to the alongshore mean of that observed bathymetry (Fig. 1). The location of \( x = 0 \) m is where the observed mean depth becomes \( h = 0 \) m, i.e., the mean shoreline. Onshore of \( x \approx 0 \) m, the model bathymetry becomes flat with

\(^1\)The Newtonian damping form used here differs slightly from that in Kennedy et al. (2000) where \( F_{br} = (1/2)(h + \eta)^{-1} \nabla \cdot \left[ \nu_{br}(h + \eta) \{ \nabla u + (\nabla u)^T \} \right] \) was used.
$h = 0.2$ m for an additional 92 m. The last 80 m of the flat region is a sponge layer (Fig. 1) that absorbs any wave energy not yet dissipated by wave breaking. At $x = -290$ m, the (observed and model) depth is $h = 7$ m, and farther offshore the model depth is flat. At the offshore end of the model domain, a second 70 m long sponge layer (Fig. 1) absorbs outgoing wave energy so that it is not reflected.

Random directionally-spread waves are generated, following, by oscillating the sea surface $\eta$ on an offshore source strip Wei et al. (1999), with 40-m cross-shore width located at $x = -445$ m in $h = 7$ m depth (light shaded region in Fig. 14). Within this strip, $\eta$ is forced at 701 individual frequencies from $0.0626$–$0.2$ Hz with 21 directional components at each frequency. The 701 frequencies and and 21 directions were sufficient so that the source standing wave problem discussed in Johnson and Pattiaratchi (2006) did not occur here. The angles (or alongshore wavenumbers) for the directional components are chosen to satisfy alongshore periodicity (Wei et al. 1999). The directional magnitudes are Gaussian so that at each frequency the mean wave angle is zero and the wave spread is constant with frequency. The forcing amplitudes (and thus the incident spectrum) are set with random phases so that the model reproduces a the wave spectra mean wave angle $\bar{\theta}$, and directional spread $\sigma_\theta$ (Appendix A) at the most offshore instrument. At the peak frequency $f_p = 0.08$ Hz, $kh = 0.45$ (where $k$ is the wavenumber) and at the highest forced frequency $f = 0.2$ Hz, $kh = 1.3$ which is within the valid Nwogu (1993) Boussinesq equation $kh$ range for wave phase speed (Gobbi et al. 2000). At the wavemaker, $H_s/h = 0.07$, and thus wavemaker nonlinearities are small.

The non-dimensional drag coefficient $c_d = 2 \times 10^{-3}$ is consistent with surfzone alongshore momentum balances (e.g., Feddersen et al. 1998) and with previous alongshore current studies using Boussinesq models (Chen et al. 2003). Biharmonic friction is necessary to dampen nonlinear aliasing instabilities in the model and the hyperviscosity is set to $\nu_{bi} = 0.3$ m$^4$ s$^{-1}$. Biharmonic friction has negligible influence on scales larger than 10 m (e.g., with $L = 10$ m, the biharmonic Reynolds number $UL^3/\nu_{bi} = 6000$). As an example, a snapshot of instantaneous $\eta$ and vorticity is shown in Fig. 2. Note that the directionally spread wave field results in wave crests with finite length (Fig. 2a).

After the model reached a statistically steady state (1000 s into the model run), 2000 model surfzone drifters were released uniformly distributed within $-240 < x < 0$ m and advected by the model’s horizontal velocities (at the reference depth $z_r = -0.531h$). Similar to the real drifters (Schmidt et al. 2003), the model drifters do not “surf” onshore at the passage of a bore. Furthermore, model drifters do not feel bore-induced turbulence and thus disperse differently than a tracer (Feddersen 2007). The model drifters were tracked for approximately 2000 s with positions output every 0.5 s. When model drifters advect onshore of $x = 0$ m, the drifter track is omitted from the dispersion calculations. Note that aside from setting the wavemaker forcing amplitudes to reproduce the most offshore ADV wave spectra, no other tuning of model coefficients has been performed to optimize the model fit to data.

### 4. Lagrangian Drifter Statistics Background

The notation, theory, techniques used to calculate the single- and two-particle Lagrangian statistics from drifter trajectories, whether observed or modeled, are introduced in this section. Note that the adjectives “single-particle” and “absolute” are synonymous when describing dispersion or diffusivities, as are the adjectives “two-particle” and “relative”.

Fig. 2. Model output snapshot, 2200 s post model spin-up, of (a) the sea surface elevation $\eta$, and (b) the vorticity for an incoming wave spread of $\sigma_\theta = 14^\circ$. The shoreline is located near $x = 0$ m. Notice significant vorticity mostly onshore of $x = -150$ m. The vertical dashed lines indicate the “inner surfzone” region ($-90 < x < -20$ m) where Lagrangian drifter analysis is performed.

**a. Single-particle statistics**

The position of the $i^{th}$ particle at time $t$ is

$$X^{(i)}(t) = X_0^{(i)} + \int_0^t v^{(i)}(\tau) \, d\tau,$$

where $v^{(i)}(\tau)$ is the Lagrangian particle velocity along its trajectory and $X_0^{(i)} = X^{(i)}(t = 0)$ is the initial particle location. The particle displacement from its original position is then

$$a^{(i)}(t|X) = X^{(i)}(t) - X_0^{(i)} = \int_0^t v^{(i)}(\tau) \, d\tau.$$
The notation “\(t|X\)” is used to indicate that the particle is at (in practice in the bin) \(X\) at time \(t\). This dependency on position is necessary as Lagrangian statistics depend on position in inhomogeneous flows. Trajectories are “binned” by final particle position when calculating absolute dispersion in inhomogeneous flows (Davis 1991). The mean displacement after time \(t\) is

\[
\bar{a}(t|X) = \langle a(t|X) \rangle
\]  

with the expectation \(\langle \cdot \rangle\) operating over all length \(t\) particle displacements that end in the bin \(X\). Note that this expectation, as long as the particle never leaves the bin \(X\), can be constructed for a single particle (not just an ensemble of many) because the time \(t = 0\) is arbitrary. Thus for a 5 second particle track there are 5 non-overlapping (but not necessarily independent) 1-second displacements. The mean displacement (5) is also the integral of the mean Lagrangian velocity

\[
\bar{a}(t|X) = \int_0^t \bar{v}(\tau|X) \, d\tau,
\]

which naturally leads to the displacement anomaly

\[
a'(t|X) = a(t|X) - \bar{a}(t|X) = \int_0^t v'(\tau|X) \, d\tau,
\]

where \(v'(\tau|X)\) is the anomalous Lagrangian velocities. The superscript “(i)” denoting particle number on \(a\) and \(a'\) is dropped as it is no longer relevant to single-particle statistics.

Quantities which depend on the “absolute” displacement \(a'\) will be designated with a superscript “(a)”. The probability density function (pdf) of displacement anomalies is \(P(a)(X,a',X',t)\) and the spreading rate is the absolute diffusivity \(\kappa^{(a)}(X,t)\). In homogeneous turbulent fluids, \(P(a)\) is expected to be Gaussian. In fluids with homogeneous turbulent statistics, the “absolute” dispersion tensor is defined as

\[
[D^{(a)}_{ij}(t)]^2 = \langle a'_i(t)a'_j(t) \rangle,
\]

and the absolute diffusivity is

\[
\kappa^{(a)}_{ij} = \frac{1}{2} \frac{d}{dt} \left[ D^{(a)}_{ij}(t) \right]^2,
\]

(Taylor 1921). However, the surfzone and nearshore (and many other oceanographic regions) do not have homogeneous turbulent velocity statistics. For example, the region just seaward of the surfzone was observed to have smaller diffusivities than within the surfzone (Spydell et al. 2007).

The absolute dispersion concepts introduced by Taylor (1921) have been extended by Davis (1987, 1991) to flows with inhomogeneous statistics such as the surfzone. In these situations, the spatially variable absolute diffusivity tensor is (Davis 1991)

\[
\kappa^{(a)}_{ij}(X,t) = \int_0^t C_{ij}(X,\tau) \, d\tau,
\]

where \(C_{ij}(X,\tau)\) is the Lagrangian velocity auto-covariance function defined as

\[
C_{ij}(X,\tau) = \langle v'_i(t|X)v'_j(-\tau+t|X) \rangle.
\]
Thus, $C_{ij}(X, \tau)$ is the $\tau$ separated autocorrelation of particle velocities binned according to the final location $X$ of the particles. The absolute dispersion is then

$$D^{(\alpha)}_{ij}(X, t) = \left[ 2 \int_0^t \kappa^{(\alpha)}_{ij}(X, \tau) \, d\tau \right]^{1/2}$$

and is a measure of the size of the ensemble average patch.

Single-particle dispersion is typically divided into two time regimes, the “ballistic” and “Brownian” regimes, which essentially assume a monotonically decaying and integrable $C_{ij}(\tau)$. For times less than the Lagrangian decorrelation time-scale, called the “ballistic” regime,

$$\left[ D^{(\alpha)}_{ij}(t) \right]^2 \sim 2E_{ij}t^2 \quad \text{for} \quad t < T_{L,ij},$$

where

$$E_{ij} = \frac{1}{2} C_{ij}(0) = \frac{1}{2} \langle v'_i(0)v'_j(0) \rangle,$$

is the Lagrangian energy. In the “Brownian” regime, times larger than the Lagrangian time-scale,

$$\left[ D^{(\alpha)}_{ij}(t) \right]^2 \sim 2\kappa^{(\alpha)\infty} t \quad \text{for} \quad t > T_{L,ij}.$$ (8)

Thus $[D^{(\alpha)}]^2$ initially scales like $t^2$ and subsequently like $t$ once the Lagrangian velocities are uncorrelated. Note that all quantities ($E_{ij}$, $T_{L,ij}$, and $\kappa^{(\alpha)\infty}$) in the above asymptotic formulae depend on position $X$.

The absolute diffusivity $\kappa^{(\alpha)}$ parameterizes eddy-fluxes for the evolution of the ensemble averaged tracer $\bar{c}(x, t)$. In an inhomogeneous flow field, $\bar{c}(x, t)$ is governed by (Davis 1987)

$$\frac{\partial}{\partial t} \bar{c} + \bar{u} \cdot \nabla \bar{c} = \nabla \cdot \left[ \int_0^t \kappa^{(\alpha)}(x, t') \cdot \nabla \bar{c}(t - t') \, dt' \right],$$ (10)

where $\bar{u}$ is the mean fluid velocity and $\kappa^{(\alpha)}$ is the particle-derived absolute diffusivity from (7). For times longer than the Lagrangian decorrelation time $T_L$ and without mean flow, tracer evolution takes the familiar form

$$\frac{\partial}{\partial t} \bar{c} = \nabla \cdot \left[ \kappa^{(\alpha)\infty}(x) \cdot \nabla \bar{c} \right].$$ (11)

One of the purposes of single-particle statistics is to estimate the diffusivities in (10) and (11).

**b. Two-particle statistics**

The separation between two particles is

$$R^{(ij)}(t) = X^{(i)}(t) - X^{(j)}(t),$$

where $X^{(i)}(t)$ and $X^{(j)}(t)$ are the locations of two distinct particles. The amount the two particles have separated from their original separation $R^{(ij)}(t = 0) = R^{(ij)}_0$ is

$$r^{(ij)}(t) = R^{(ij)}(t) - R^{(ij)}_0$$
and the separation anomaly is
\[ r^{(ij)}(t|X_m, R_0) = r^{(ij)}(t) - \bar{r}(t|X_m, R_0) \]
where \( \bar{r}(t|X_m, R_0) = \langle r^{(ij)}(t|X_m, R_0) \rangle \) and ensemble averages are taken over all particle pairs with the same initial separation \( R_0 \) and same initial location of the pair \( X_m \) – the pair’s initial midpoint. Thus unlike the notation used for single-particle statistics, “\( t|X_m, R_0 \)” for two-particles means at \( t = 0 \) the midpoint of the particles is \( X_m \) and the initial separation is \( R_0 \). The probability density function of particle separation anomalies is denoted by \( P^{(r)}(X_m, R_0, r', t) \), where the superscript \( r \) denotes “relative”.

The “width” of this pdf is given by the relative dispersion
\[ D^{(r)}(X_m, R_0, t) = \left[ \langle r_i'(t|X_m, R_0)r_j'(t|X_m, R_0) \rangle \right]^{1/2} \]
which indicates how far particles have separated from their initial separation. The relative diffusivity \( \kappa^{(r)}_{ij} \) is the rate two-particles separate and is defined as
\[ \kappa^{(r)}_{ij}(X_m, R_0, t) = \frac{1}{2} \frac{d}{dt} \left[ D^{(r)}_{ij}(X_m, R_0, t) \right]^2. \]

In 2D homogeneous isotropic turbulence, the statistics of particle separations are known. In the inverse-energy cascade range (the inertial subrange), i.e., for length-scales larger than the injection scale of the turbulence, the velocity wavenumber spectrum scales as \( E(k) \sim k^{-5/3} \) and the relative dispersion scalings are
\[ P^{(r)}(r') \sim \exp(-|r'|^{2/3}) \]  
(12a)
\[ \left[ D^{(r)}(t) \right]^2 \sim t^3 \]  
(12b)
\[ \kappa^{(r)} \sim \left[ D^{(r)} \right]^{4/3} \]  
(12c)
These scalings (12) are called Richardson’s Laws and have been observed in simulations of isotropic inertial subrange 3D turbulence (Boffetta and Sokolov 2002a), simulations of 2D turbulence (Boffetta and Sokolov 2002b), and laboratory experiments of 2D turbulence (Jullien et al. 1999). Although (12b,c) are justified from dimensional arguments (Obukhov 1941a,b; Batchelor 1950), (12a) cannot be theoretically derived but is rather obtained by analogy with diffusion (Richardson 1926).

In the 2D turbulent enstrophy cascade, length-scales smaller than the injection scale of the turbulence where \( E(k) \sim k^{-3} \), the relative dispersion and diffusivity scale as (Lin 1972)
\[ \left[ D^{(r)}(t) \right]^2 \sim \exp(t) \]  
(13)
\[ \kappa^{(r)} \sim \left[ D^{(r)} \right]^2. \]
These scalings (13), have been recently observed in laboratory experiments of 2D turbulence (Jullien 2003). At long times, when particles separations are much larger than the largest eddies, the particles move independently and the relative dispersion asymptotes to absolute dispersion, i.e.,
\[ \frac{1}{2} [D^{(r)}]^2 \to [D^{(a)}]^2, \]  
(14)
and both scale as \( \sim t^1 \). For random spatially and temporally correlated velocity fields (i.e., at the smallest separations), it is straightforward to show that \( [D^{(r)}]^2 \sim t^2 \) and \( \kappa^{(r)} \sim D^{(r)} \).
Fig. 3. The observed (squares) and modeled (lines) Eulerian wave statistics versus cross-shore coordinate $x$: (a) significant wave height $H_s$, (b) the significant wave height to total water depth $d$ ratio $\gamma = H_s/d$, (c) directional wave spread $\sigma_\theta$, and (d) the observed water depth $h$. The “inner-surfzone” region ($-90 < x < -20$) is shaded. In (b) the total water depth $d$ is mean depth $h$ plus the mean setup $\bar{\eta}$. In addition, the inner-surfzone dashed lines indicate the previously observed (mean ± standard deviation) $\gamma$ range (Raubenheimer et al. 1996).

Unlike single-particle statistics whose primary utility is quantifying eddy-fluxes for the ensemble averaged tracer evolution, two-particle statistics are primarily useful for determining the structure of the flow field (turbulent or not). However, similar to the ensemble averaged patch, the spreading of the “typical” tracer patch (i.e., $P^{(r)}$), can be modeled by a diffusion-like equation with two-particle statistics quantifying the patch spreading (see Richardson 1926; Kraichnan 1966; Spydell et al. 2007).

The concept of 2D turbulence and the associated Lagrangian relative dispersion are based upon a constant depth fluid. Depth variation will affect the vorticity dynamics central to 2D turbulence, and lead to a non-isotropic, non-homogeneous turbulence. Thus the surfzone eddy field will not strictly follow canonical 2D turbulence. However, in general the surfzone bottom slope (here 0.025) is small, and these 2D turbulence concepts (both Eulerian and Lagrangian signatures) will be compared to the observed and modeled 2-particle statistics.

5. Model / Data Comparison: Eulerian Statistics

Prior to comparing observed and modeled Lagrangian statistics, an Eulerian wave and current comparison is performed. Wave statistics were observed at 3 ADV locations. The observed and modeled cross-shore variation of the significant wave height $H_s$ compare well (Fig. 3a), however the ADV locations are
not ideal for a model test. Offshore $H_s = 0.5$ m and increases in shallower depths (Fig. 3d) due to shoaling until the about $x \approx -130$ m where $H_s$ decreases. At the innermost ADV ($x = -107$ m), located in the outer-surfzone, the model overpredicts $H_s$. There were no ADV observations in the inner-surfzone ($-90 < x < -20$ m). Within the inner-surfzone, modeled $\gamma$, the ratio of significant wave height to total water depth, gently increases (Fig. 3b) and is consistent with a $\gamma$ parameterization (dashed curves in Fig. 3b), based upon extensive field observations, which depends upon local beach slope and $kh$ (Raubenheimer et al. 1996). The waves are normally incident at all ADVs as the observed mean wave angle magnitude $|\bar{\theta}| < 2^\circ$ is within the ADV orientation error. The modeled waves are also normally incident with $|\bar{\theta}| < 1^\circ$ at all cross-shore locations. The wave maker input was chosen so that the modeled and observed $\sigma_\theta$ at the most offshore ADV were in approximate agreement. In the inner-surfzone, modeled $\sigma_\theta$ increases (as observed in Herbers et al. 1999) and may be due surfzone eddies refracting waves analogous to the increasing $\sigma_\theta$ due to shear-wave induced wave refraction (Henderson et al. 2006). In addition to bulk (sea-swell integrated)
moments (i.e., $H_s$), modeled and observed sea-surface elevation spectra at the ADVs are in good agreement in the sea-swell band (not shown).

Observed and modeled drifter velocities are spatially binned and averaged to obtain Eulerian mean and fluctuating velocity statistics (Fig. 4). Modeled binned statistics are essentially alongshore uniform. The observed and modeled drifter derived mean currents are weak (typically $<0.1$ m s$^{-1}$, red arrows in Fig. 4). Within the inner surfzone at the $x = -60$ m bin, the alongshore averaged mean alongshore current is $0.005$ m s$^{-1}$ which is not significantly different from zero. Neither observed or modeled mean cross-shore currents show any indication of long-lived rip currents. Similar to the model results of Johnson and Pattiaratchi (2006), short-lived ($\sim 100$ s) episodic rip currents occurred in the model. The observed and modeled drifter derived mean currents are weak (typically $<0.1$ m s$^{-1}$). Within the inner surfzone at the $x = -60$ m bin, the alongshore averaged mean alongshore current is $0.005$ m s$^{-1}$ which is not significantly different from zero. Neither observed or modeled mean cross-shore currents show any indication of long-lived rip currents. Similar to the model results of Johnson and Pattiaratchi (2006), short-lived ($\sim 100$ s) episodic rip currents occurred in the model. The observed and modeled drifter derived mean currents are weak (typically $<0.1$ m s$^{-1}$). Within the inner surfzone at the $x = -60$ m bin, the alongshore averaged mean alongshore current is $0.005$ m s$^{-1}$ which is not significantly different from zero. Neither observed or modeled mean cross-shore currents show any indication of long-lived rip currents. Similar to the model results of Johnson and Pattiaratchi (2006), short-lived ($\sim 100$ s) episodic rip currents occurred in the model. The observed and modeled drifter derived mean currents are weak (typically $<0.1$ m s$^{-1}$).

The similarity between the observed and modeled bulk wave statistics up to the outer surfzone (Fig. 3a,c), the inner-surfzone $\gamma$ (Fig. 3b), and binned drifter velocities (Fig. 4) indicate that the Boussinesq model is reasonably representing surfzone processes. This is a pre-requisite to a Lagrangian drifter dispersion model-data comparison. However, the Eulerian data set is limited and more detailed Boussinesq model-data comparison with more extensive Eulerian field data sets will be performed.


a. Single Particle (Absolute) Lagrangian Statistics

The modeled and observed drifter trajectories are used to calculate $P^{(r,a)}(D^{(r,a)})^2$, $\kappa^{(r,a)}$, as described in Section 4 and compared with each other. The absolute displacement pdf $P^{(a)}(X, a'_x, a'_y, t)$, with $X$ the inner-surfzone bin ($-90 < x < -20$ m), was calculated from both the observed and modeled displacements (Fig. 5). As discussed in Spydell et al. (2007), for $t \leq 16$ s, the observed $P^{(a)}$ is polarized in the cross-shore direction $a'_x$ (1st and 2nd columns of Fig. 5) and alongshore $a'_y$ polarized for $t \geq 64$ s (last column of Fig. 5). Thus, on average, a delta function tracer release initially spreads more quickly in the cross-shore, becomes roughly circular at $t \approx 16$ s, and subsequently elongates more rapidly in the alongshore. The observed and modeled $P^{(a)}$ are similar for all $t$ (compare top and bottom panels of Fig. 5). At short times ($t = 1, 4$ s) the observed $P^{(a)}$ is broader in $y$, and has more outliers in both $x$ and $y$ than modeled due to GPS position errors.

To further compare modeled and observed displacement pdfs, one-dimensional (1D) displacement pdfs are defined as

$$P^{(a)}(a'_x) = \int_{-\infty}^{\infty} P^{(a)}(a'_x, a'_y) da'_y$$

(15)

(and similarly for alongshore displacements, $P^{(a)}(a'_y)$). These 1D pdfs are then non-dimensionalized by their respective observed dispersion $D^{(a)}_{ij}$ so that pdf shapes at different times can be directly compared. Both observed and modeled non-dimensional pdfs, for both along- and cross-shore displacements, approximately
collapse to a Gaussian curve as expected for a 2D random flow field (Fig. 6). For the shortest time \( t = 1 \) s, the modeled \( \bar{P}(a_x, a_y, t) \) is skewed toward \( +a_x \), due to steep waves (large \( +u \) velocities) inducing large onshore \( +a_x \) displacements (Fig. 6a). GPS position error likely obscures this in the observations. For \( t \geq 16 \) s, both modeled and observed \( \bar{P}(a_x, t) \) have negatively skewed longer-than-Gaussian tails (i.e., at \( > 2|a_x|/D_{xx}^{(o)} \), Fig. 6a). This is because in the inner-surfzone bin neither modeled nor observed drifter can have large \( +a_x \) displacements as drifters would end up beached, but large offshore \( -a_x \) displacements are possible. The non-dimensional alongshore displacement pdfs \( \bar{P}(a_x, t) \) also deviate somewhat from Gaussian with the modeled \( \bar{P}(a_x, t) \) having longer than Gaussian tails (Fig. 6b). The observed \( \bar{P}(a_x, t) \) is negatively skewed due to poor sampling while the modeled \( \bar{P}(a_x, t) \) is symmetric for all \( t \) as expected for a alongshore uniform surfzone and normally incident waves.

Observed and modeled single-particle (absolute) cross- \( (D_{xx}^{(o)}) \) and alongshore \( (D_{yy}^{(o)}) \) dispersion are calculated for the inner-surfzone bin as described in Section 4. The observed and modeled \( [D_{xx}^{(o)}]^2 \) are similar (Fig. 7a). At \( t < 5 \) s, observed and modeled \( [D_{xx}^{(o)}]^2 \) have similar power laws (between 1–2). For \( 20 < t < 300 \) s, the power-law for the model becomes more ballistic \( (t^1 \) power law) while the one for the observations remains relatively unchanged from \( t < 20 \) s. Both observed and modeled \( [D_{yy}^{(o)}]^2 \) are Brownian \( (t^1 \) power law) for \( t > 300 \) s. The modeled and observed \( [D_{xx}^{(o)}]^2 \) have similar magnitude for \( t < 100 \) s, and at later times \( (200 < t < 1000 \) s\) the modeled is 1.5 to 3 times larger than observed (observed \( [D_{xx}^{(o)}]^2 \) is between 400-800 m² over this time). The observed and modeled \( [D_{yy}^{(o)}]^2 \) are also similar (Fig. 7b). Both modeled and observed \( [D_{yy}^{(o)}]^2 \) show ballistic \( (20 < t < 300 \) s) and Brownian regimes \( (t > 600 \) s). For \( t < 200 \) s the observed \( [D_{yy}^{(o)}]^2 \) is larger than the modeled, possibly due to GPS errors. Thereafter \( (200 < t < 1000 \) s\), the modeled \( [D_{yy}^{(o)}]^2 \) is 1 to 2 times larger than the observed (observed \( [D_{yy}^{(o)}]^2 \) is between 1300-3000 m².) The modeled and observed absolute diffusivities, \( \kappa_{xx}^{(o)} \) and \( \kappa_{yy}^{(o)} \) (7) for
Fig. 6. Observed (dots) and modeled (curves) pdf $\tilde{P}^{(a)}$ of absolute (a) $x$ displacements $a'_x$, and (b) $y$ displacements $a'_y$ at times $t = 1, 4, 16, 64, 128, 256$ s (colors from blue to orange). Both the pdfs and displacements are scaled by the standard deviation of the displacements ($D^{(a)}_{xx}$ and $D^{(a)}_{yy}$, respectively) at that time. Only times out to $t = 256$ s, are shown to minimize sampling error in the observed $\tilde{P}^{(a)}$. The Gaussian (solid black lines) distribution is indicated.

use in (10), are also similar (Fig. 8), particularly at shorter times ($t < 60$ s, Fig. 8a,b). At longer times, the modeled $\kappa_{xx}^{(a)}$ is larger than the observed and the modeled $\kappa_{yy}^{(a)}$ becomes approximately $2.5 \times$ larger than the observed (Fig. 8c,d). Model diffusivity estimates seaward of the surfzone (see Spydell et al. 2007) are not discussed.

b. Two Particle (Relative) Lagrangian Statistics

One-dimensional separation pdfs $\tilde{D}^{(r)}(r'_x)$ and $\tilde{P}^{(r)}(r'_y)$ are defined similarly to the one-dimensional
Fig. 7. Observed (dashed) and modeled (solid) single particle dispersion (a) $[D_{xx}^{(a)}]^2$ and (b) $[D_{yy}^{(a)}]^2$ versus time $t$. Both $t^1$ and $t^2$ scalings are indicated as thin lines.

$\bar{P}^{(a)}$ (15). As previously found for the observations (Spydell et al. 2007), the modeled non-dimensional separation pdfs $\bar{P}^{(c)}$ for small initial separations, $|R_0| < 4$ m, follow Richardson scaling (12a) for all times (only $t < 256$ s is shown in Fig. 9 where there are a sufficient number of observations for quality comparison). However, the modeled separation pdfs become more Gaussian for larger $|R_0|$ and longer times (not shown). As discussed in Spydell et al. (2007), the Richardson scaled $\bar{P}^{(c)}$ (as opposed to Gaussian) imply that drifter pairs do not move independently due to a self similar interacting eddy field over a range
of length-scales. That the observed and modeled non-dimensional $\bar{P}^{(r)}$ agree well for both cross- and alongshore separations (Fig. 9a and b, respectively), indicates that observed and modeled surfzone eddy field separating drifters are similar.

The observed and modeled cross- ($[D_{xx}^{(r)}]^2$) and alongshore ($[D_{yy}^{(r)}]^2$) relative dispersion are calculated for the inner-surfzone region from drifter separations as described in Section 4b. Unlike the absolute dispersion, the modeled $[D_{xx}^{(r)}]^2$ is less than the observed for all times (Fig. 10a,b). The largest differences occur for small times and small $[D_{xx}^{(r)}]^2$ ($< 10$ m$^2$) in part due to GPS position errors ($\pm 1$ m). The modeled $[D_{xx}^{(r)}]^2$ grows slowly on wave period time-scales ($t < 20$ s) which is also seen, albeit less strongly, in the observations (Fig. 10a). For times $t > 300$ s, both observed and modeled $[D_{xx}^{(r)}]^2$ and $[D_{yy}^{(r)}]^2$ have greater than $t^2$ power-law dependence, although this strong time dependence is not as clear in the observations. This indicates the
Fig. 9. Observed (dots) and modeled (curves) pdf $P^{(r)}(\bar{r})$ of relative (a) $x$ separations $r_x'$, and (b) $y$ separations $r_y'$ at times $t = 1, 4, 16, 64, 128, 256$ s (colors from blue to orange). Both the pdfs and displacements are scaled by the standard deviation of separations ($D_{xx}^{(r)}$ and $D_{yy}^{(r)}$, respectively) at that time. Only times out to 256 s, are shown to minimize sampling error in the observed $P^{(r)}$. Richardson ($\sim \exp(-|r'|^2/3)$, dashed black lines) and Gaussian (solid black lines) distributions are indicated.

Examination of the modeled relative diffusivity dependence upon the relative dispersion shows that $\kappa_{xx}^{(r)} \sim [D_{xx}^{(r)}]^2$ for $10 < D_{xx}^{(r)} < 20$ m and $\kappa_{yy}^{(r)} \sim [D_{yy}^{(r)}]^2$ for $5 < D_{yy}^{(r)} < 25$ m (gray thick lines in Fig. 10c,d), indicating enstrophy cascade scaling (13). At these length-scales, $[D^{(r)}]^2$ should grow exponentially (gray thick lines in Fig. 10a,b). However, detecting this is difficult as it occurs for such a small range of length-scales. For length-scales smaller than the onset of enstrophy cascade scaling ($D^{(r)} \leq 5$ m), modeled and
Fig. 10. The relative dispersion (a) \( [D^{(r)}_{xx}]^2 \) and (b) \( [D^{(r)}_{yy}]^2 \) versus time \( t \) and The relative diffusivity (c) \( \kappa^{(r)}_{xx} \) and (d) \( \kappa^{(r)}_{yy} \) versus separation \( D^{(r)} \) (\( D^{(r)}_{xx} \) and \( D^{(r)}_{yy} \), respectively) for initial separations \( |R_0| < 4 \) m. Observations are dashed lines and model results are solid lines. Power-law scaling are indicated as thin solid lines. The enstrophy cascade scalings (13) are indicated as thick gray lines.

Observed \( \kappa^{(r)}_{yy} \sim [D^{(r)}_{yy}]^{1} \) as expected for purely random but correlated velocity fields. Both modeled \( \kappa^{(r)}_{xx} \) and \( \kappa^{(r)}_{yy} \) are weakly scale dependent at length-scales between \( 25 < D^{(r)} < 40 \) m but become at least \( \kappa^{(r)} \sim [D^{(r)}]^1 \) for length-scales larger than 40 m. At these larger length-scales the \( [D^{(r)}]^2 \) appears to scale \( \sim t^3 \) for \( t > 1000 \) in both \( x \) and \( y \), an indicator of a 2D turbulent inverse-energy cascade.

Overall, modeled and observed two-particle statistics are comparable. In particular, the observed and modeled separation pdf shapes are very similar (Fig. 9). This combined with the similar power law scalings for the relative dispersion and the scale-dependent diffusivities (despite quantitative disagreement), indicate that in both the model and observations there is a turbulent eddy field with a range of length scales. The similarity between modeled and observed Lagrangian statistics motivates using the model to investigate the
Fig. 11. Modeled alongshore velocity spectra $G_{vv}(f)$ versus frequency $f$ for the full model $u$, irrotational $u_\phi$, and rotational $u_\psi$ velocities (see legend) at $x = -60$ m. The features of the cross-shore velocity spectra are similar.

underlying processes driving surfzone Lagrangian dispersion.

7. Velocity Decomposed Dispersion

Any two-dimensional velocity field can be written as the sum of an irrotational velocity due to velocity potential $\phi$ and the rotational velocity due to the curl of the streamfunction $\psi$, i.e.,

$$u = \nabla \phi + \nabla \times \psi$$  \hspace{1cm} (16)

where $\nabla$ is the two-dimensional gradient operator. The irrotational velocity $u_\phi = \nabla \phi$ is the divergent part of the flow and the rotational velocity $u_\psi = \nabla \times \psi$ can have non-zero vorticity. For the $\sigma_\theta = 14^\circ$ model run, the full model velocity field was output every $\Delta t = 0.5$ s for 5000 s after model spin-up. From this, $\phi$ and $\psi$ are calculated at each time step by solving the elliptic equations

$$\nabla^2 \phi = \nabla \cdot u, \quad \text{and} \quad \nabla^2 \psi = \zeta,$$  \hspace{1cm} (17)

where the vorticity $\zeta = \nabla \times u$. The alongshore boundary conditions for both $\phi$ and $\psi$ are periodic. At the onshore boundary, $\psi = 0$ and $\partial_x \phi = 0$. At the offshore boundary, $\psi = \int \langle v \rangle dx$ and $\phi = \int \langle u \rangle dx$ where $\langle \rangle$ represents an alongshore average. Both boundaries are within the sponge layer. From $\phi$ and $\psi$, $u_\phi$ and $u_\psi$ are estimated. The resulting root-mean-square (rms) errors (averaged in time, the alongshore, and over the region where drifters were released, $-240 < x < 0$ m) from the velocity decomposition are small ($\text{rms}||u - (u_\phi + u_\psi)|| < 0.01$ m s$^{-1}$).
At sea-swell frequencies (0.05 < f < 0.3 Hz), the velocities are largely irrotational. The $u$ and $u_\psi$ spectra are nearly identical for both cross- and alongshore (Fig. 11) components and are 2 or more orders of magnitude larger than the $u_\psi$ spectra. The $u_\phi$ velocities are dominated by variability at sea-swell frequencies. The modeled rotational motions are dominated by low frequencies and the $u_\psi$ spectra is red. At low ($f < 0.005$ Hz), the full $u$ and $u_\psi$ spectra are similar (gray and thick curves in Fig. 11) and are larger than the $u_\phi$ spectra. A snapshot of the modeled vorticity is shown in Fig. 2b.

Drifters are seeded into and advected with the full $u$, irrotational $u_\phi$, and rotation $u_\psi$ velocity fields. Examples of 1000-s long drifter tracks from the three velocity fields are shown in Fig. 12. Overall, the $u_\psi$ advected drifter tracks are smooth, have large 50–100 m displacements (Fig. 12b), and visually appear as low-passed full $u$ drifter tracks (Fig. 12a). The $u_\phi$ advected drifter tracks have much smaller displacements and are dominated by oscillations induced by high frequency surface gravity wave (Fig. 12c). These oscillations are also observed in the full $u$ drifter displacements (Fig. 12a). Note that the sum of the $u_\phi$ advected and $u_\psi$ advected drifter trajectories do not and should not equal the full model $u$ trajectories.

Inner-surfzone absolute $D^{(o)}$ and relative $D^{(r)}$ drifter dispersions are calculated for each of the 3 velocity fields (Fig. 13). Results for $[D^{(r)}]^2$ are shown for the $u$, $u_\phi$, and $u_\psi$ velocity fields. The results are qualitatively similar for the decomposed absolute dispersion $[D^{(o)}]^2$. At short times $t < 10$ s, the cross-shore $[D_{xx}^{(x)}]^2$ dispersion for the full $u$ is nearly identical to the $u_\phi$ dispersion (blue and red curves in Fig. 13a), resulting from random surface gravity waves. This is consistent with the $u_\phi$ spectra dominant at $f > 0.05$ Hz. However, at longer time-scales $t > 100$ s, the $u_\psi$-induced $[D_{xx}^{(r)}]^2$ asymptotes to the full $u$ dispersion (blue and green curves in Fig. 13a), where the $u_\phi$-induced $[D_{xx}^{(r)}]^2$ is 2 orders of magnitude smaller. The full $u$ dispersion $[D_{yy}^{(r)}]^2$ follows the $u_\psi$ dispersion for all times $t > 10$ s (Fig. 13b). These results demonstrate that surfzone dispersion is dominated by rotational motions and the $[D^{(r)}]^2$ power-law time-dependence suggests
The asymptotic diffusivity depends only on two quantities, the Lagrangian energy $E_{ij}$ and time-scale...
As the dispersion for \( t > O(10) \) s is dominated by rotational velocities, the \( E_{ij} \) used is derived only from rotational velocities, i.e.,

\[
E_{ij}^{(\psi)} = \frac{1}{2} \langle v_i^{(\psi)} v_j^{(\psi)} \rangle. \tag{18}
\]

The irrotational surface gravity wave contribution to \( E_{ij} \) is not included because although its zero-lag Lagrangian velocity covariance (e.g., \( E_{ij}^{(\phi)} \)) is substantial, irrotational motions do not contribute to the long-time dispersion. Thus, using the full \( E_{ij} \), and known Lagrangian time-scale, results in asymptotic diffusivity predictions that are too large (since \( \kappa_{ij}^{(a)\infty} = E_{ij} T_{L,ij}/2 \)). The Lagrangian time-scale is then calculated from (9b) using the particle derived \( \kappa_{ij}^{(a)\infty} \) (see Fig. 8) and the \( E_{ij}^{(\psi)} \), i.e.,

\[
T_{L,ij} = \frac{\kappa_{ij}^{(a)\infty}}{2 E_{ij}^{(\psi)}}.
\]

Examining only the diagonal components, with \( \kappa_{ij}^{(a)\infty} = [0.75, 4.00] \) m²s⁻¹ and \( E_{ij}^{(\psi)} = [0.005, 0.006] \) m²s⁻² yields \( T_L = [75, 333] \) s, considerably longer than \( T_L = [7, 54] \) s for the day one observations (Spydell et al. 2007). This discrepancy results from the including irrotational velocities in \( E_{ij} \) used to calculate the observed \( T_L \). With the \( u_\psi \) derived \( T_L \) and \( E_{ij}^{(\psi)} \), both the ballistic and Brownian regimes for the modeled \( |D^{(a)}| \) are well predicted (see dashed \( t \) and \( t^2 \) lines in Fig. 13), except for \( |D_{xx}^{(a)}| \) for \( t < 10 \) s which is surface gravity wave dominated. This further demonstrates the dominance of vorticity (rotational motions) in absolute as well as relative dispersion.

8. Surfzone Eddies, Vorticity Variability, and Directional Wave Spread

As shown in Section 7, dispersion is dominated by rotational (vorticity) motions (i.e., surfzone eddies) rather than irrotational ones. The mechanism by which surfzone eddies are generated and how drifter dispersion is influenced is now addressed. In general, surfzone eddies have many possible generation mechanisms. Shear waves generate surfzone vorticity variability (Oltman-Shay et al. 1989) which in numerical models can spin up into eddies (e.g., Allen et al. 1996). However, shear waves require significant mean alongshore current shear (e.g., Bowen and Holman 1989) which was not present for the normally incident waves on this day of observations. Alongshore bathymetric variability may also play a role in generating surfzone eddies. Very low frequency (f < 0.004 Hz) rotational motions were observed to be coupled to a rip channel morphology (MacMahan et al. 2004). However, spatially and temporally variable radiation stress forcing (i.e., wave-groups, which is essentially the wave-averaged result of a random directionally spread wave field) were required to properly model the underlying very low frequency variability (Reniers et al. 2007). Similar to the modeling results of Johnson and Pattiaratchi (2006), here a rich surfzone rotational velocity field (e.g., Fig. 2b) is generated on an alongshore uniform bathymetry. As discussed in Peregrine (1998), alongshore gradients in breaking wave heights act as a vorticity source in shallow water dynamics.

The effect of alongshore non-uniform wave breaking on vorticity is seen by taking the curl of (3) (neglecting higher order terms) which results in

\[
\frac{\partial \zeta}{\partial t} + \ldots = \nabla \times F_{br} \tag{19}
\]

where \( \ldots \) represents the standard vorticity advective and stretching terms. The curl of the dispersive (\( \nabla \times F_d = O((kh)^2) \)), bottom stress, and biharmonic friction terms in (3) are neglected. To see how this term

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acts as a vorticity source, consider normally incident waves with alongshore varying amplitude. As these waves enter the surfzone, depth-limited breaking only occurs where the waves are largest thus resulting in finite crest-length broken waves and non-zero $F_{br}$ (see schematic in Fig. 14). In this case, $F_{br}$ is cross-shore $(x)$ oriented and varies in the alongshore direction, thus $\nabla \times F_{br}$ is non-zero, generating vorticity.

On alongshore uniform bathymetry, alongshore variable wave amplitude and thus finite breaking crest lengths are the result of directionally spread wave fields. The larger $\sigma_\theta$ the shorter the average breaking crest length. Surfzone vorticity, and hence Lagrangian dispersion, should then depend upon the incident wave directional spread $\sigma_{\theta_0}$. To test this idea, four additional model simulations with identical wave conditions except for the incident $\sigma_{\theta_0}$ were performed, resulting in 5 total runs to be analyzed with $\sigma_{\theta_0} = 0^\circ, 4^\circ, 7^\circ, 14^\circ, 20^\circ$.

The $\sigma_{\theta_0} = 0^\circ$ simulation is not realistic for a surfzone as there was zero alongshore velocity at all time due to the alongshore uniformity. No real beach has perfect alongshore uniform bathymetry and wave field. However, the $\sigma_{\theta_0} = 0^\circ$ run is interesting as an idealized example of the limit of infinite-crest length breaking waves. This simulation clearly resulted in no vorticity generation and negligible (single- and two-particle) drifter dispersion. Results from this simulation are thus not shown.

The cross-shore dependence of the wave spread $\sigma_\theta$ is similar for each of the different $\sigma_{\theta_0} = 4^\circ, 7^\circ, 14^\circ$ and $20^\circ$ simulations (Fig. 15a). For each $\sigma_{\theta_0}$, the wave spread decreases as the surfzone is approached and then, for all but the $\sigma_{\theta_0} = 20^\circ$ simulation, increases through the inner-surfzone until the shore is reached. Recall that for the $\sigma_{\theta_0} = 14^\circ$ run, the modeled $\sigma_\theta(x)$ matched the observations (see also Fig. 3c). The mean vorticity for all $\sigma_{\theta_0}$ is zero at all $x$ as bores can only generate (potential-) vorticity anomalies (Bühler 2000). However, the model vorticity standard deviation std($\zeta$) (based on a time- and alongshore average) increases with larger incident $\sigma_{\theta_0}$ (Fig. 15b), and increases within the inner-surfzone where the majority of wave dissipation occurs. Similarly, Kennedy (2005) showed that increasing $\sigma_\theta$ increases the magnitude of the fluctuating rotational velocities. Well offshore of the surfzone ($x < -200 \text{ m}$) the vorticity variability is...
small as few eddies generated in the surfzone were able to propagate that far offshore.

Inner-surfzone vorticity wavenumber spectra $G_{ζζ}(k_ζ)$ are constructed for all $σ_θ_0$ by averaging the instantaneous vorticity alongshore wavenumber periodogram over the inner-surfzone ($−90 < x < −20$ m) and over $t$. For all $σ_θ_0$, $G_{ζζ}$ is red and spread over a large range of alongshore wavenumbers $k_ζ$ (Fig. 16). For $k_ζ > 5 \times 10^{-3}$ cpm (cycles per meter), $G_{ζζ}$ is larger for increasing $σ_θ_0$. Thus the increased vorticity variability with larger $σ_θ_0$ (Fig. 15) is spread over 10–200 m length-scales. For all $σ_θ_0$, $G_{ζζ}$ appears to have two differing wavenumber dependencies. Specifically, for the $σ_θ_0 = 14^°$ and $20^°$ runs, $G_{ζζ}$ falls off very rapidly for $k_ζ > 0.05$ cpm and more gently for $0.01 < k_ζ < 0.05$ cpm, i.e., at the transitional wavenumber $k_ζ \approx 0.05$ cpm, the $G_{ζζ}$ power-law dependence changes. The two other $σ_θ_0$ runs also show two $G_{ζζ}(k_ζ)$ regimes, but with a smaller transitional wavenumber for decreasing $σ_θ_0$, consistent with longer breaking crest-lengths injecting energy at longer length-scales.

The increased vorticity variability induced by increasing $σ_θ_0$ also results in larger inner-surfzone relative dispersion (Fig. 17). For $t < 10$ s, the cross-shore dispersion $[D_{xx}^{(r)}]^2$ is similar for all $σ_θ_0$ (Fig. 17a) as these time-scales are too short for vorticity motions to separate drifters. At longer times $t > 10$ s, $[D_{xx}^{(r)}]^2$ is larger with increasing $σ_θ_0$ as surfzone eddies separate the drifters. In addition, for larger $σ_θ_0$, significant cross-shore drifter separation $([D_{xx}^{(r)}]^2)$ begins at earlier times as the increased vorticity variance (Fig. 15b) at smaller length-scales increases (Fig. 16). At $t = 2000$ s, order of magnitude differences in $[D_{xx}^{(r)}]^2$ exist for the various $σ_θ_0$. For example, with $σ_θ_0 = 20^°$ and $4^°$, drifters have cross-shore separated an average of $D_{xx}^{(r)} = 45$ m and $D_{xx}^{(r)} = 33$ m, respectively (Fig. 17a). In general, $[D_{yy}^{(r)}]^2$ is larger for increased $σ_θ_0$ (Fig. 17b). The $[D_{yy}^{(r)}]^2$ power-law scaling is similar for all $σ_θ_0$, only the magnitude varies. At times $10 < t < 1000$ s, the $σ_θ_0 = 20^°$ $[D_{yy}^{(r)}]^2$ is slightly larger than for $σ_θ_0 = 14^°$, whereas for $t > 1000$ s they

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**Fig. 15.** Modeled (a) wave directional spread $σ_θ$ and (b) vorticity standard deviation std($ζ$) versus $x$ for different incident directional spreads $σ_θ_0$. The open squares in (a) are the ADV observations.
are the same. Furthermore, the $\sigma_{\theta_0} = 14^\circ$ and $20^\circ$ runs have largely similar $G_{\zeta\zeta}(k_y)$ (Fig. 16), possibly indicating a vorticity saturation. The $\sigma_{\theta_0}$ dependence of absolute dispersion is qualitatively similar to that of relative dispersion (not shown).

The modeled relative dispersion indicates the presence of a 2D turbulent enstrophy and inverse-energy cascade with an injection length-scale of approximately $20 \text{ m}$ for $\sigma_{\theta_0} = 14^\circ$. Although the dispersive velocities are rotational and the vorticity is spread over wavenumber space, the 2D turbulent character of the modeled surfzone remains to be quantified from the Eulerian model data. An Eulerian statistic useful for classifying enstrophy and inverse-energy cascade regions is the velocity structure function $S_v(\Delta y)$ defined as

$$S_v(\Delta y) = \langle [v(y + \Delta y) - v(y)]^2 \rangle,$$

where $v$ is the instantaneous alongshore velocity and the average $\langle \cdot \rangle$ is over time and the alongshore direction $y$. In 2D turbulence, dimensional arguments (e.g., Kellay and Goldburg 2002), lead to

$$S_v(\Delta y) \sim (\epsilon \Delta y)^{2/3}, \quad L_0 > \Delta y > y_{in}$$

$$S_v(\Delta y) \sim \beta^{2/3}(\Delta y)^2, \quad \Delta y < y_{in}$$

for an inverse-energy and enstrophy cascade, respectively. In (21), $L_0$ is the largest length-scale where velocities are correlated, $y_{in}$ is the injection length-scale at which the 2D turbulence field is forced, $\beta$ is the enstrophy injection rate and $\epsilon$ is the energy injection rate. The structure function scalings (21) are analogous to the $E \sim k_y^{-3}$ and $E \sim k_y^{-5/3}$ wavenumber spectra scalings for enstrophy and inverse-energy cascade regions, respectively. Note however that an $S_v \sim (\Delta y)^2$ scaling is also possible at small $\Delta y$ for purely
random but spatially correlated velocities. The structure function $S_v(\Delta y)$ (20) is readily calculated for the different $\sigma_{\theta_0}$ at various $x$ locations from the Boussinesq model Eulerian data (Fig. 18).

For the $\sigma_{\theta_0} = 14^\circ$ run, $S_v(\Delta y)$ suggests an enstrophy and inverse-energy cascade (Fig. 18c). Specifically, for inner-surfzone locations (e.g., $x = -34$ m), the structure function follows $S_v \sim (\Delta y)^2$ for approximately $\Delta y < 10$ m (steepest dashed line in Fig. 18c) suggesting the presence of an enstrophy cascade region. At larger scales, approximately $20 < \Delta y < 150$ m, $S_v \sim (\Delta y)^{2/3}$ indicating the presence of an inverse-energy cascade region (gently sloping dashed line in Fig. 18c). The transition length-scale $y_{\text{in}}$ from enstrophy to inverse-energy cascades is between 10–20 m, consistent with the transition scale inferred from the relative dispersion statistics (Section 6b). At $\Delta y > L_0 = 200$ m, $S_v(\Delta y)$ approaches a constant as alongshore velocities at these separations are uncorrelated. Seaward of the surfzone (e.g., $x = -189$ m, thin lines in Fig. 18c), no inverse-energy cascade ($S_v \sim (\Delta y)^{2/3}$) region is observed, as there is no breaking-
wave vorticity injection. The other $\sigma_{\theta_0}$ runs exhibit similar behavior, but with weaker overall $S_v$ and with $y_{in}$ and $L_0$ increasing for decreasing $\sigma_{\theta_0}$. This is consistent with the larger breaking crest-lengths injecting vorticity at larger length-scales. The $\sigma_{\theta_0} = 4^\circ$ run is unique in that there is no inverse-energy cascade region at any cross-shore locations.

From both Eulerian (structure function) and Lagrangian (relative dispersion) analysis, the modeled surfzone appears 2D turbulent-like. The turbulence magnitude (rms vorticity or relative dispersion) and the length-scales $L_0$ and $y_{in}$ depend upon $\sigma_{\theta_0}$ ($y_{in} \approx 10 - 20$ m for the $\sigma_{\theta_0} = 14^\circ$ simulation). For larger $\sigma_{\theta_0}$,
the modeled inner-surfzone structure function follows inverse-energy cascade scaling (i.e., $S_v \sim (\Delta y)^{2/3}$) for $\Delta y > 20$ m (Fig. 18). This is consistent with the modeled relative dispersion as the non-dimensionalized separation pdfs follow the Richardson scaling and $[D^{(r)}]^2 \sim t^3$ for $D^{(r)} > 20$ m (Fig. 17). Furthermore, the length-scales for the enstrophy cascade are similar for the Eulerian ($\Delta y < 10$ m) and Lagrangian ($5 < D^{(r)} < 20$ m) analyses.

9. Summary

Surfzone drifter dispersion was observed on a beach with small normally incident directionally spread waves (Spydell et al. 2007). For these conditions, surfzone Lagrangian drifter dispersion was simulated with a time-dependent wave-resolving Boussinesq model. The limited observed Eulerian (wave properties, mean currents) statistics are well reproduced by the model. The model reproduces the observed absolute dispersion statistics with approximately Gaussian displacement pdfs and comparable along- and cross-shore dispersions (and diffusivities). The long-time model alongshore absolute diffusivities are $2.5 \times$ larger than observed. The observed relative dispersion is reasonably well reproduced by the model. Both observed and modeled non-dimensionalized separation pdfs are Richardson-like. The modeled $[D^{(r)}_{xx}]^2$ and $[D^{(r)}_{yy}]^2$ are smaller than observed. For short times, this is likely in part due to GPS error in the observations. The modeled and observed relative dispersions have approximately the same power law time dependence (stronger than $[D^{(r)}]^2 \sim t^1$) as well as the relative diffusivity having the same power law scale dependence $(\kappa^{(r)} \sim [D^{(r)}]^n$ with $1 \leq n \leq 2$). Both enstrophy and inverse-energy cascade regions are identified in the modeled relative dispersion.

The model velocity field was decomposed into irrotational (surface gravity waves) and rotational (vorticity) motions. Higher frequency ($f > 0.01$ Hz) motions are dominated by irrotational velocities, surface gravity waves, and lower frequency ones ($f < 0.005$ Hz) by rotational velocities. Drifters are advected within the irrotational and rotational velocity fields. At times longer than $t \approx 30$ s, absolute and relative drifter dispersion are dominated by the rotational velocity field indicating the importance of surfzone eddies (vorticity) in drifter dispersion. Alongshore gradients in breaking wave dissipation generate vorticity (e.g., Peregrine 1998) over a range of scales. On an alongshore uniform beach, a directionally spread wave field is required for finite breaking crest-lengths. Simulations with increased incident $\sigma_\theta_0$ result in increased rms vorticity over a broader range of length-scales, giving rise to increased drifter dispersion. The velocity structure function $S_v(\Delta y)$ generally shows regions with both enstrophy and inverse-energy cascade scalings. For larger $\sigma_\theta_0$, $S_v(\Delta y)$ magnitude increases and both the upper ($L_0$) and lower ($y_{\text{in}}$, the injection scale of the turbulence) length-scale limits of the inverse-energy cascade decreases Both Eulerian ($S_v(\Delta y)$) and Lagrangian (two-particle) statistics reveal that the modeled surfzone is a quasi 2D-turbulent fluid. For the $\sigma_\theta_0 = 14^\circ$, these Eulerian and Lagrangian statistics generally indicate an enstrophy cascade (approximately 5–10 m length-scales) and inverse-energy cascade (approximately 20–100 m length-scales).

Acknowledgments. This research was supported in part by CA Sea Grant and ONR. Sea Grant support was through the National Sea Grant College Program of the U.S. Department of Commerce’s National Oceanic and Atmospheric Administration under NOAA Grant NA04OAR4170038, project #01-C-N, through the California Sea Grant College Program and the California State Resources Agency. The observations were collected in collaboration with R. T. Guza and W. E. Schmidt. Staff from the Integrative
Oceanography Division of SIO, designed and built the drifters, and were instrumental in acquiring the field observations. Discussions with David Clark, R. T. Guza, and Steve Henderson provided valuable insight. The Boussinesq model funwaveC was developed by F. Feddersen and is freely available as open-source software at http://iod.ucsd.edu/~falk/models.html

APPENDIX A Definition of Wave Statistics

The frequency directional sea-surface elevation spectrum is given by $E_{\eta \eta}(f, \theta)$ where $f$ is the frequency and $\theta$ the wave direction. The frequency spectrum $G_{\eta \eta}(f)$ is the integral over all directions,

$$G_{\eta \eta}(f) = \int_{-\pi}^{\pi} E_{\eta \eta}(f, \theta) d\theta,$$

so that

$$\text{Var}(\eta) = \int_{0}^{\infty} G_{\eta \eta}(f) df$$

and the significant wave height $H_s$ is defined as

$$H_s = 4 \left[ \int_{SS} G_{\eta \eta}(f) df \right]^{1/2},$$

where the integral is over the sea-swell band (SS) of 0.05–0.3 Hz. The bulk (sea-swell band frequency-integrated) wave angle $\bar{\theta}$ and directional spread $\sigma_\theta$ are defined as (Kuik et al. 1988)

$$\bar{\theta} = \arctan \left[ \frac{\int_{SS} \int_{-\pi}^{\pi} \sin(\theta) E_{\eta \eta}(f, \theta) d\theta df}{\int_{SS} \int_{-\pi}^{\pi} \cos(\theta) E_{\eta \eta}(f, \theta) d\theta df} \right]$$

and

$$\sigma_\theta^2 = \frac{\int_{SS} \int_{-\pi}^{\pi} \sin^2(\theta - \bar{\theta}) E_{\eta \eta}(f, \theta) df d\theta}{\int_{SS} \int_{-\pi}^{\pi} E_{\eta \eta}(f, \theta) df d\theta}$$

Direct estimates of the directional spectrum are not required to calculate $\bar{\theta}$ and $\sigma_\theta$. Instead both are functions of the the lowest bulk Fourier directional moments $a_2$ and $b_2$ (Kuik et al. 1988) as described in Herbers et al. (1999) which depend upon the the $u$ and $v$ cross-spectra. For the field data these wave statistics are estimated from the $u$ and $v$ spectra converted with linear theory to sea surface elevation spectra and for the model output, the $\eta$, $u$ and $v$ spectra are similarly used.

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Draft: April 17, 2008