

# The initiation and evolution of transverse dunes: A literature review

Paul Jarvis

## 1 Introduction

## 2 Fluid flow

The first consideration when considering sedimentary processes is the driving flow. Although it is possible for sediment transport to occur in laminar flow (Charru and Mouilleron-Arnould, 2002; Charru et al., 2004; Mouilleron et al., 2009), the majority of research into the dynamics of ripples and dunes has considered turbulent conditions (Andreotti et al., 2002; Kroy et al., 2002; Langlois and Valance, 2007; Wierschem et al., 2008; Franklin and Charru, 2011; Charru and Franklin, 2012; Andreotti et al., 2012; Franklin, 2015), a reflection of the fact that these are relevant to the majority of natural systems. Therefore, we here summarise the fundamentals of turbulent flow, before discussing descriptions of wall bounded turbulent flow, both over flat and wavy surfaces.

### 2.1 Turbulence

The flow of any incompressible Newtonian fluid is described by the Navier-Stokes (NS) equation and the incompressibility relation:

$$\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} = -\frac{\nabla p}{\rho} + \nu \nabla^2 \mathbf{u}; \quad (1)$$

and

$$\nabla \cdot \mathbf{u}. \quad (2)$$

These describe conservation of momentum and mass respectively where  $\mathbf{u}(\mathbf{x}, t)$  is the velocity field,  $p(\mathbf{x}, t)$  the pressure field,  $\rho$  the fluid density and  $\nu$  the kinematic viscosity

of the fluid, whilst  $\mathbf{x}$  and  $t$  are the spatial and temporal variables. If we define the components of the stress tensor  $\boldsymbol{\tau}$  to be

$$\tau_{ij} = \rho\nu \left( \frac{\partial u_i}{\partial x_j} + \frac{\partial u_j}{\partial x_i} \right), \quad (3)$$

then the NS equation can be re-written as

$$\rho \left( \frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} \right) = -\nabla p + \nabla \cdot \boldsymbol{\tau}. \quad (4)$$

It is useful to define the quantity vorticity,  $\boldsymbol{\omega} = \nabla \times \mathbf{u}$  which provides a measure of the local angular momentum of fluid elements. Vorticity is a useful quantity for a number of reasons. Firstly, it can be shown to obey a simpler equation than  $\mathbf{u}$ ,

$$\frac{\partial \boldsymbol{\omega}}{\partial t} + (\mathbf{u} \cdot \nabla) \boldsymbol{\omega} = (\boldsymbol{\omega} \cdot \nabla) \mathbf{u} + \nu \nabla^2 \boldsymbol{\omega}. \quad (5)$$

Secondly, it can be shown for two-dimensional flows that vorticity must be conserved within flow interiors; it cannot be created or destroyed but only advected or diffused. Vorticity can only be created at the boundaries of fluid domains, before being transferring into the interior. The origin of this creation is the strong velocity gradients within boundary layers.

The third use of vorticity is that it provides a starting point for a possible definition of turbulence. Turbulence is a poorly defined quantity. However, it is possible to picture turbulent eddies as regions of vorticity, with turbulence being an evolving field of eddies (Davidson, 2004). Here we will use the following definition of turbulence from Corrsin (1961), which has been modified by Davidson (2004): “Incompressible fluid turbulence is a spatially complex distribution of vorticity which advects itself in a chaotic manner in accordance with (5). The vorticity field is random in both space and time, and exhibits a wide and continuous distribution of length and time scales”.

With a definition of turbulence, we can now proceed with a description of turbulent flow. This is achieved by decomposing  $\mathbf{u}$ ,  $p$  and  $\boldsymbol{\tau}$  into mean and fluctuation fields e.g.  $\mathbf{u} = \bar{\mathbf{u}} + \mathbf{u}'$  where the bar denotes the mean field and the prime the fluctuation. The time-averaged NS equation is then shown to be given by

$$\rho \left( \frac{\partial \bar{u}_i}{\partial t} + \bar{u}_j \partial_j \bar{u}_i \right) = -\partial_i \bar{p} + \partial_i (\bar{\tau}_{ij} + \tau_{ij}^R) \quad (6)$$

where the Reynolds stresses  $\tau_{ij}^R = -\rho \bar{u'_i u'_j}$ . Additionally, time-averaging the incompressibility relation yields both  $\partial_i \bar{u}_i = 0$  and  $\partial_i u'_i = 0$ . Reynolds stresses are not stresses

in a physical sense. Physical stresses correspond to momentum fluxes between different fluid particles. Reynolds stresses correspond to momentum flux between the mean and fluctuation fields.

With this definition of Reynolds stresses, it is possible to manipulate the equations of motion to obtain an equation for them. This is algebraically arduous and will not be repeated here but the end result contains a term  $\partial_k(u'_i \bar{u}'_j u'_k)$ . The triple correlations here are unknown, but deriving an equation for them leads to a further set of unknowns  $u'_i \bar{u}'_j u'_k u'_l$ . The governing equations for these involves fifth-order correlations and so on. Hence, this statistical description of turbulence is undetermined - there are more unknowns than equations. This is called the closure problem. Numerous empirical models have been developed for a closure relation which completes the system of equations. The most widely used are eddy viscosity models, based on the principle that dissipation of momentum from the mean flow to the fluctuations by the Reynolds stresses is analogous to dissipation of momentum by viscous stresses. Hence, an analogous eddy viscosity  $\nu_t$  can be defined as (Boussinesq, 1887)

$$\tau_{ij}^R = \rho \nu_t (\partial_j \bar{u}_i + \partial_i \bar{u}_j) - \frac{\rho u'_k \bar{u}'_k \delta_{ij}}{3}. \quad (7)$$

How  $\nu_t$  is determined depends on the geometry of the flow. Since we are considering wall-bounded two-dimensional flows the most common method is to use Prandtl's mixing length model (Prandtl, 1925) (an analogue to the kinetic theory of viscosity) which gives

$$\nu_t = l_m^2 |\partial_z \bar{u}_x|, \quad (8)$$

where  $l_m$  is the mixing length which can be determined empirically. Whilst the model relies on dubious assumptions, it has been shown to be very succesful for one-dimensional shear flows.

## 2.2 Wall-bounded flow

We will initially consider turbulent flow over a flat, smooth, solid surface, with the mean flow directed in the x-direction, and the z-direction denoting the upward normal to the surface. The mean velocity in the interior of the flow (outside the boundary layer) is  $\bar{\mathbf{u}} = (U_0, 0, 0)$  and the flow is unbounded as  $z \rightarrow \infty$ . The boundary layer thickness  $\delta$  can be defined as the value of  $z$  at which  $\bar{u}_x(z) = 0.99U_0$ , although other measures exist (Pope, 2000). A significant amount of work has gone into studying the structure of this boundary layer. It has been postulated (Prandtl, 1925), and since verified both experimentally (Wei and Willmarth, 1989) and numerically (Kim et al., 1987), that there exists an inner layer close to the wall ( $y/\delta \ll 1$ ) where  $\bar{u}(x)$  depends on the viscous scales and is independent

of  $\delta$  and  $U_0$ . The viscous lengthscale is defined as

$$\delta_\nu = \frac{\nu}{u_\tau}, \quad (9)$$

where  $u_\tau$  is the friction velocity, defined as

$$u_\tau = \left( \frac{\tau_w}{\rho} \right)^{1/2}, \quad (10)$$

and  $\tau_w$  is the shear stress exerted on the wall. This postulate can be used to find the velocity profile within the inner layer. Defining  $y^+ = y/\delta_\nu$ , it can be shown that for small  $y^+$  there exists a viscous sublayer where (Pope, 2000)

$$\bar{u}_x(z) = \frac{u_\tau z}{\delta_\nu}, \quad (11)$$

and for  $y^+ \gg 1$ , there is a so called log-law region (von Kármán, 1930) where

$$\bar{u}_x(z) = \frac{u_\tau}{\kappa} \ln \left( \frac{z}{z_0} \right). \quad (12)$$

Here  $\kappa$  is the von Kármán constant and  $z$  is a constant of integration called the hydrodynamical roughness, found by matching to the profile in the viscous sublayer. Experimental and numerical studies have shown that the viscous sublayer typically extends to  $y^+ < 5$  and the log-law region for  $y^+ > 30$ ,  $y/\delta < 0.3$  (Kim et al., 1987; Wei and Willmarth, 1989), with  $\kappa \approx 0.4$ . The transition between these two parts of the inner layer is called the buffer layer. Beyond the inner layer, there exists an outer layer, which has predominantly been described empirically (Pope, 2000).

The mixing length for wall-bounded flow is given by (Pope, 2000)

$$l_m = \kappa z \left[ 1 - \exp \left( -\frac{z}{\alpha \delta_\nu} \right) \right], \quad (13)$$

where  $\alpha \approx 25$  is the van Driest number.

For a rough surface, the roughness can be described by a bed of grains of characteristic diameter  $d$  fixed at  $z = 0$  (Charu et al., 2013). For  $d/\delta_\nu \lesssim 5$  the effect of the roughness in the log-law region is negligible, as the disturbance to the flow is contained within the viscous sublayer. However, at larger  $d$ , the flow is said to be hydrodynamically rough, and experimental observations are used to estimate the hydrodynamic roughness (Bagnold, 1941; Kamphuis, 1974). Various works use  $l_m = \kappa(z + z_0)$  as the mixing length for a hydrodynamically rough flow (Fourrière et al., 2010).

## 2.3 Flow over a perturbed profile

We now consider the flow over a sinusoidal profile  $\zeta = \zeta_0 \cos(kx)$ , where  $\zeta_0$  is the amplitude of the perturbations to the surface, and  $k$  is the wavenumber. If the gradient of the profile is small ( $k\zeta_0 \ll 1$ ), then the resultant perturbation to the flow will be linear. For example, the resultant perturbation to the shear-stress in the  $x$ -direction at the fluid surface can be expressed as (Kroy et al., 2002; Charru et al., 2013)

$$\tau_b = \frac{\hat{\tau}e^{ikx} + \hat{\tau}^*e^{-ikx}}{2}, \quad \text{where} \quad \hat{\tau} = \tau_w(A + iB)k\zeta_0. \quad (14)$$

The periodic functional form of  $\tau_b$  is due to the periodic nature of the solid boundary. The contribution from  $A$  is in-phase with the surface disturbance and that from  $B$  is in quadrature.

Depending on the relative values of the different length scales in the problem ( $k^{-1}$ ,  $\delta_\nu$ ), the perturbation to the flow occurs in different hydrodynamic regimes. For the case that the flow disturbance is confined to the viscous sublayer ( $k\delta_\nu \ll 1$ ), the flow disturbance due to the surface perturbations is laminar, and has a two-layered structure (Charru et al., 2013): an inner viscous layer, and an outer inertial layer. Remember, both of these layers are within the viscous sublayer and are different from the much larger inner and outer layers that describe the coarser structure of the general boundary layer. The shear-stress on the solid surface can be evaluated, and both  $A$  and  $B$  in equation (14) are found to both be positive. This means the maximum of the shear stress is always located upstream of the crests (Benjamin, 1958; Charru and Hinch, 2000; Charru et al., 2013).

The other possibility is that the flow disturbance extends beyond the viscous sublayer ( $k\delta_\nu > 1$ ). It can again be shown that the flow disturbance has a multi-layered structure, but here the perturbation is turbulent. Jackson and Hunt (1975) and Sykes (1980) considered two-layers: a thin, turbulent layer outside the viscous sublayer dominated by gradients in Reynolds stresses, overlain by an inertia-dominated layer that continued up to the outer layer of the boundary layer. This description was later improved with the inclusion of a middle layer which improved matching (Hunt et al., 1988; Belcher and Hunt, 1998). As for small peaks, the maximum shear stress is again found upstream of the crests (Hunt et al., 1988; Belcher and Hunt, 1998; Kroy et al., 2002; Charru et al., 2013).

As well as the laminar and turbulent regime, there is a transitional regime where neither viscous effects nor turbulent fluctuations can be ignored. Thus far, only empirically verified simple parameterisations of the viscous sublayer thickness have been used in describing this regime (Abrams and Hanratty, 1985; Frederick and Hanratty, 1988), and a true physical understanding of the behaviour is lacking.

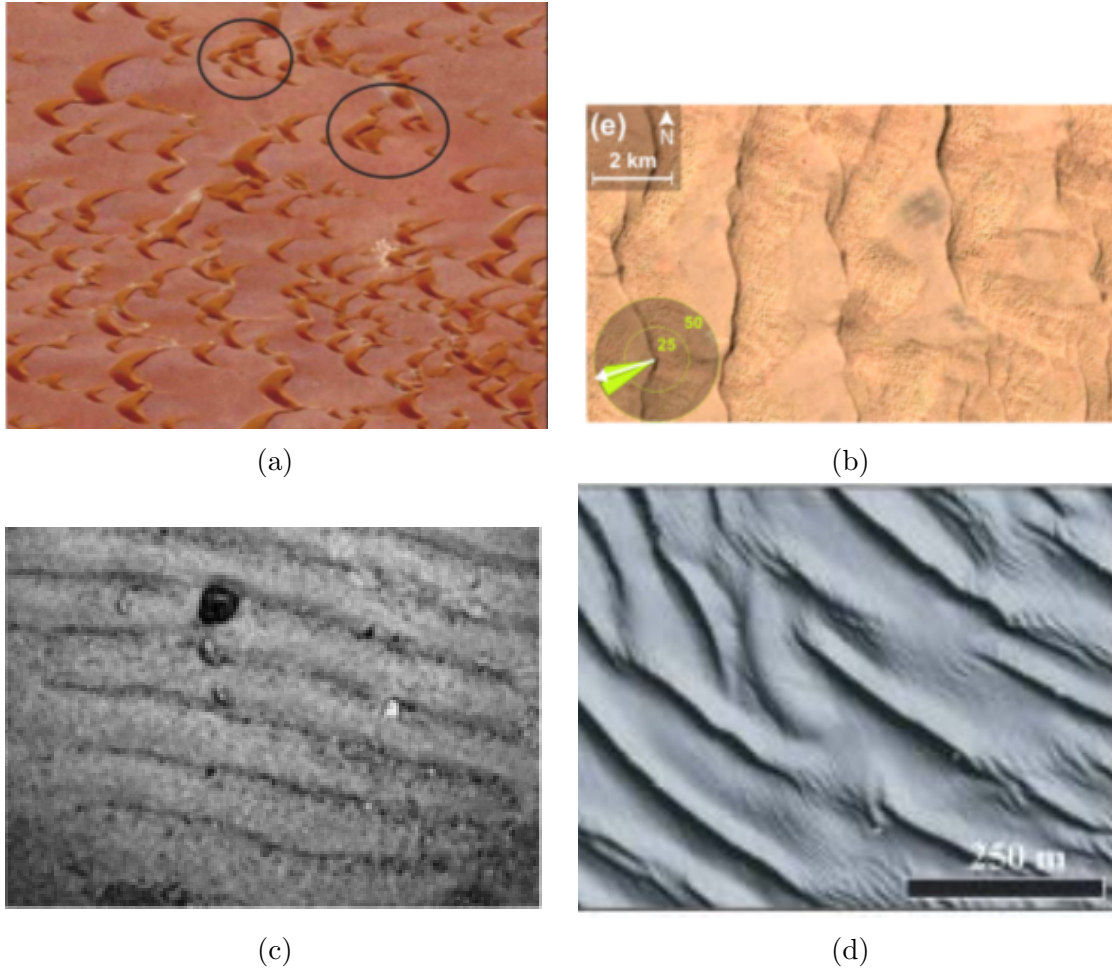


Figure 1: Examples of bedforms. a) A field of barchan dunes from Morocco. (Durán et al., 2011). b) Transverse dune field in the East Taklamakan desert, China (Gao et al., 2015). c) Subaqueous ripples in Zion National Park, USA (Andreotti et al., 2012). d) Large subaqueous dunes in San Francisco Bay, USA (Barnard et al., 2006).

### 3 Types of sedimentary bedforms

There are a variety of different types of sedimentary bedforms found in nature, with the diversity reflecting the hugely variable conditions under which erosion and sedimentation processes occur. Most generally, there are two types of environment in which bedforms can be found: subaerial; and subaqueous. Subaerial bedforms result from the flow of air over a granular medium, such as a desert, beach or a volcanic ash deposit. Subaqueous bedforms are created by the action of water, and can be found on river-beds, shallow coastal regions, or deeper continental shelves. Within this hierarchy, further variability in shape and size of bedform originates from spatial and temporal variations in flow-speed, flow-depth, particle size, particle density and local topography. Examples of different types of bedforms are shown in figure 1.

## 4 Modes of sediment transport

In order for an overlying fluid to displace any sediment, it is necessary for the shear stress exerted on the bed surface to be sufficient to lift a particle up and over its neighbours. Hence, the lift force provided by the flow must be sufficient to overcome gravity. This balance is quantified through the Shields number

$$\Theta = \frac{\tau}{(\rho_p - \rho_f)gd}, \quad (15)$$

where  $\tau$  is the surface shear stress,  $\rho_p$  and  $\rho_f$  are the densities of the particles and fluid respectively,  $g = 9.81 \text{ m s}^{-2}$  is the acceleration due to gravity and  $d$  is the particle diameter. Hence, there is some critical value of the Shields number at which motion becomes possible.

Broadly, there are three mechanisms describing how sediment can be transported by an overlying fluid (Bagnold, 1941). For Shields numbers just above the threshold for motion, bedload transport (reptation) will occur, whereby particles slide and roll across the surface. With increasing flow velocity, these particles can be lifted higher into the flow, and ‘hop’ across the bed following ballistic trajectories, a process called saltation. Once the size of turbulent fluctuations in the vertical fluid velocity are comparable to the settling velocity of the particles, they can become suspended in the fluid and are transported as suspended load. For the purpose of the rest of this review, all sediment transport can be assumed to occur through reptation or saltation unless otherwise specified.

There is a finite limit to the amount of sediment a given flow can transport. This is quantified through the saturated sediment flux  $q_{\text{sat}}$  (Bagnold, 1941; Owen, 1964). Therefore, there must also be a lengthscale  $L_{\text{sat}}$  and/or timescale  $T_{\text{sat}}$  over which this saturation can be achieved. Expressions for  $q_{\text{sat}}$  scales depend on the mechanics of transport but a general partial differential equation for the sediment flux  $q$  can be given as (Charru et al., 2013)

$$T_{\text{sat}} \frac{\partial q(x, t)}{\partial t} + L_{\text{sat}} \frac{\partial q(x, t)}{\partial x} = q_{\text{sat}} - q(x, t). \quad (16)$$

This equation is valid for a unidirectional flow.

## 5 Initiation of transverse structures

## 6 Coarsening

## 7 Open questions

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