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Green’s Law Revisited: Tidal Long-Wave Propagation in Channels With Strong Topography

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INTRODUCTION

Long-wave propagation in narrow channels of variable width and depth has been the subject of analysis for more than 150 years. One of the first major results was that of Green [1837], after whom Green’s functions were named. Green’s Law states that the wave amplitude \( \zeta \) and tidal transport \( Q \) vary with width \( b \) and depth \( h \) according to \( \zeta \approx b^{-1/2}h^{1/4} \) and \( Q \approx b^{-1/2}h^{1/4} \). This solution is of limited utility because it is restricted to inviscid, infinitesimal waves in channels with no mean flow and weak topography (those with topographic scale \( L \gg \lambda \), wavelength \( \lambda \)). An analytical perturbation model including finite-amplitude effects, river flow, and tidal flats has been used to show that (1) wave behavior to lowest order is a function of only two nondimensional parameters representing, respectively, the strength of friction at the bed and the rate of topographic convergence/divergence; (2) two different wave equations with nearly constant coefficients can be derived that together cover most physically relevant values of these parameters, even very strong topography; (3) a single, incident wave in a strongly convergent or divergent geometry may mimic a standing wave by having a \( \pm 90^\circ \) phase difference between \( Q \) and \( \zeta \) and a very large phase speed, without the presence of a reflected wave; (4) channels with strong friction and/or strong topography \( (L \ll \lambda) \) show very large deviations from Green’s Law; and (5) these deviations arise because both frictional damping and the direct dependence of \( |Q| \) and \( \zeta \) on topography (topographic funneling) must be considered.

long-wave propagation in channels is nonlinear and has not been solved exactly, but two approximate solutions are derived herein. Both include river flow, finite amplitude, friction, and topography. Together, these solutions are adequate to represent wave propagation through narrow channels with arbitrarily strong (but smooth) depth and width variation. One, the “standard solution,” is a generalization of Green’s solution. It has asymptotically constant coefficients in two circumstances. Either friction and acceleration effects on the wave number are both weak relative to topography (strong topography), or friction and topography are weak relative to acceleration effects (weak topography).

The other, the “critical solution,” represents wave propagation when the flow is dominated by bed friction. This solution becomes exact (constant coefficient) at a particular, critical convergence rate where acceleration and topographic effects are equal and of opposite sign in the wave equation. Critical convergence is also the greatest possible convergence rate for inviscid waves. Above critical convergence (in the supercritical range), viscous disturbances damp exponentially rather than propagate \([\text{Lighthill, 1978}]\). Thus friction must be included in any representation of long-wave propagation over strong topography.

Both of these solutions are derived from a perturbation expansion analysis of the long-wave equation. The strength of finite-amplitude nonlinearities (those that give rise to overtide and residual circulations and increase wave propagation speed) in this equation is scaled by the barotropic nonlinearity parameter \( \varepsilon \equiv \zeta/H \), the ratio of wave amplitude to mean depth \([\text{Stokes, 1847}]\). In the presence of bed friction, finite-amplitude effects appear not only in the convective acceleration term in the wave equation, but also in the friction term itself \([\text{Parker, 1984}]\). Although it may approach unity in very shallow systems, \( \varepsilon \) is small but finite in most systems and is thus appropriate as a perturbation expansion parameter for finite-amplitude effects. The other likely form of analysis for this problem, that based on characteristics, is not particularly appropriate, because at critical convergence the wave equation is diffusive, not hyperbolic, in the terminology of \([\text{Whitham, 1974}]\) and \([\text{LeBlond, 1978}]\).

Finite amplitude is not the only source of nonlinearity in long-wave propagation. Polynomial representation of the bed stress term \( c_D |U|U \) (where \( c_D \equiv O(3 \times 10^{-3}) \) is the drag
coefficient and $U$ is along-channel velocity) gives rise to terms as high as cubic in $U$ [Dronkers, 1964] or even quintic [Godin and Gutierrez, 1986]. River flow acts through the convective accelerations to slow movement of the incident, landward propagating wave and speed that of the reflected, seaward propagating wave in a convergent channel. It also greatly augments the linear and nonlinear effects of friction. Finally, tidal flats adjacent to a channel cause nonlinear effects through the convective acceleration term in the wave equation [Speer and Aubrey, 1985; Friedrichs and Aubrey, 1988]. Momentum is lost on the rising tide as water pours out over flats and decelerates because of strong friction. A further loss occurs on the falling tide as water decelerates to slow movement of the reflected wave alone (i.e., without a reflected wave) mimics a standing wave with incident and reflected waves of equal amplitude. The model also shows that Green’s Law is greatly modified by the presence of friction and/or strong topography. Finally, finite amplitude modifies wave celerity by a factor of $1 + (3a/2)h/b$. This reduces to the classical result of $1 + (3a/2)$ [Lamb, 1932] for $b_r$ equal to $b$.

Application of this model to actual estuarine topographies is not discussed because multiple tidal constituents and overides are vital to tidal predictions and because the intricacies of real topographies would obscure rather than illuminate the major features of wave behavior. This analysis concentrates instead on a very general analysis of the response of the fundamental wave alone to topography and friction. The overide wave (to be included in an extended version of the model) exhibits more complex behavior because of its nonlinear forcing. Nonetheless, a careful comparison has been made of the present model predictions for the geometry of the Columbia River estuary to both tidal height observations and results from a multiconstituent model [Giese and Jay, 1989] for three river flow levels and neap, spring, and mean tides. The performance of the present model for $M_2$ is systematically better than that of Giese and Jay, primarily because finite-amplitude and topographic effects on the wave number have been included here.

**Governing Equations**

**Derivation of a Wave Equation**

The first step in any analysis of wave propagation in tidal channels is to derive a one-dimensional wave equation from sectionally integrated expressions for conservation of mass and along-channel momentum. Use of sectionally integrated equations is justified because channel curvature, the ratio of width to tidal wavelength, and the Kelvin number (width divided by barotropic Rossby radius) are all small in the channels considered herein. The tidal flow is nearly two-dimensional in the along-channel $(x)$ and vertical $(z)$ directions, allowing cross-channel integration of the continuity and momentum equations. Tidal transport $(Q(x,t))$ is a “natural” variable in a stratified tidal flow in a way that the velocity $U(x, z, t)$ is not, because of the large flood/ebb changes in vertical distribution of the horizontal tidal flow that occur in many systems [Jay, 1991]. Thus a vertically integrated approach in which only $Q$ and elevation $e(x, t)$ are calculated is appropriate. The cross-sectionally integrated dimensional equations for conservation of mass and along-channel momentum are

$$
\frac{\partial Q}{\partial z} + b \frac{\partial Q}{\partial t} = 0
$$

and

$$
\frac{\partial Q}{\partial t} + \frac{1}{A} q \frac{Q^2}{\partial \xi} = 0
$$

where

- $\xi$ = along-channel distance, positive landward;
- $t$ = time;
- $A$ = channel cross-section, equal to $b(h + \xi)$;
- $g$ = acceleration due to gravity;
- $\zeta$ = bed stress divided by water density.

A circumflex indicates a dimensional variable.
Inclusion of tidal flats in a one-dimensional analysis requires that (1) and (2) be modified to account for the momentum flux to and from the tidal flats. This can be accomplished using the ideas of Speer [1984]. Speer assumed that no momentum is conveyed in the $\hat{x}$ direction by flow over these flats. Water moving over the flats on a rising tide loses its momentum to bed friction, so that upon return to the channel as the water level falls, its velocity is zero; it must then be accelerated to the ambient flow velocity. Moreover, the dissipation over the flats is assumed to be small relative to that which occurs in the channels. Clearly, these two assumptions become inconsistent if the area of tidal flats is too large. However, because velocities are typically smaller along the channel margins that feed water to the tidal flats than in mid-channel, the following analysis will allow the tidal flats at each $\hat{x}$ to be as wide as the channel itself. C. Friedrichs and O. Madsen (manuscript in preparation, 1991) have considered the effects of arbitrarily wide tidal flats in a channel of constant width and depth with strong friction. This analysis, which they refer to as “kinematic” or “zero inertia,” corresponds roughly to the critical convergence case in a channel with uniform (in $\hat{x}$) width and depth, but with wide tidal flats; $\varepsilon$ is modified by the presence of these flats.

Consider first the effects of tidal flats on conservation of mass (equation (1)). Define the time-varying total width $b_T(x, \hat{t})$ as the sum of a constant (in time) stream width $b(x)$ and a tidal flat width $\Delta b(x, \hat{t})$ that varies with the elevation of the tide (Figure 1). Intuitively, it is obvious that the along-channel divergence of $\hat{Q}$ is related to $b_T$ and not $b$, because the surface elevation $\xi$ rises over the entire width, not just the stream width. More formally, Speer, 1984, the presence of tidal flats causes an additional momentum flux $\hat{q}_T = \Delta b \partial \xi / \partial \hat{t}$. Adding the additional momentum flux and assuming a linear tidal flat profile, conservation of mass

\[
\frac{\partial Q}{\partial \hat{t}} + \frac{\partial \hat{Q}}{\partial \hat{x}} = \frac{\partial \xi}{\partial \hat{x}} + \Delta b \frac{\partial \xi}{\partial \hat{t}} + 2 \hat{\xi} \frac{\partial \xi}{\partial \hat{t}} \frac{\partial}{\partial \hat{t}} \quad (3)
\]

at tidal frequency, where $\xi^I$ is the tidal amplitude. Tidal flats have thus two effects on mass conservation. First, $b_T$ rather than $b$ appears as the multiplier of $\partial \xi / \partial \hat{t}$, and second, a new nonlinear term has appeared that is associated with the tidal cycle variation in tidal flat width. The latter term is important to overtide generation but will not influence the basic tidal motion. Friedrichs et al. [1990] have argued that, for a wide tidal flat, the form of the tidal flat profile is important because this profile changes the amount of storage over the flats. Any of several different tidal flat profiles could be incorporated in the present analysis. Use of the simplest possible (linear) profile herein minimizes algebraic complexity and is consonant with the generally poor knowledge of tidal flat profiles in most estuaries.

The effect of the momentum flux associated with flow of water on and off tidal flats in the along-channel momentum equation is the product of change in velocity $-\hat{Q}/\hat{A}$ as water moves onto the flats times the tidal flat momentum flux $\hat{q}_T$ [Speer, 1984]. With some further manipulation and taking $\Delta b / b_T \equiv \Delta b / b_T$, (2) may be written

\[
\frac{\partial \hat{Q}}{\partial \hat{t}} + 2\hat{b}_T - \hat{b} \frac{\partial \hat{Q}}{\partial \hat{x}} \hat{A} - \frac{\partial \hat{b}}{\partial \hat{t}} \frac{\partial \hat{Q}}{\partial \hat{x}} \hat{A} + \frac{\partial \hat{g} \hat{\xi}}{\partial \hat{t}} + \hat{b}_T \hat{b}_T = 0 \quad (4)
\]

Thus as one intuitively expects, the momentum flux associated with the presence of tidal flats modifies the convective acceleration terms in the motion equation.

The above two first-order differential equations in $\hat{Q}$ and $\hat{\xi}$ representing conservation of mass and momentum, (3) and (4), may be combined by cross differentiation to form a second-order wave equation in one variable. In principle, it is immaterial whether this wave equation is in $\hat{Q}$ or $\hat{\xi}$. However, the approximate form of the bed stress term to be used herein is much simpler if the wave equation is formed in $\hat{Q}$. Thus $\partial \hat{Q}/\partial \hat{t}$ of (3) is subtracted from $\partial \hat{Q}/\partial \hat{t}$ of (4), and (3) is used to express the pressure gradient in terms of $\hat{Q}$. Several bilinear pressure gradient terms are neglected because they do not affect the tidal motion before $O(\varepsilon^2)$. At the same time, two simplifications are made in the convective accelerations. First, the only tidal-frequency terms in the bilinear convective accelerations are those involving $\hat{Q}_i \hat{Q}_j / \hat{A} = \hat{U}_{ik} \hat{Q}_{jk}$, where $\hat{Q}_i$ is the river flow transport and $\hat{U}_{ik}$ is the river

\[
\begin{align*}
\text{Fig. 1. Model geometry, with MWL as mean water level, HW as high water, LW as low water, } & \Delta \hat{b}(x, \hat{t}) \text{ as time-variable tidal flat width, } \Delta b(x) \text{ as mean tidal flat width, } b \text{ as momentum conveying width, } \hat{b}_T \text{ as total width, } h \text{ as mean depth, and } \xi \text{ as tidal elevation. } \\
\text{Positive } x \text{ direction is landward.}
\end{align*}
\]

width, $h$ as mean depth, and $\xi$ as tidal elevation. The positive $x$ direction is landward.
flow velocity. River flow is assumed steady and enters the system only at the upstream end, so that $\dot{Q}_R$ is invariant with $\dot{x}$ and $t$. Second, because the convective accelerations will ultimately scale as $O(\varepsilon)$, $\dot{A} = \dot{h} \dot{b}$ rather than $\dot{A} = \dot{b}h(1 + \xi \dot{h})$ may be used in the convective accelerations. The neglected terms in $\dot{Y}_h$ are $O(\varepsilon)$ relative to the remaining terms and are thus only $O(\varepsilon^2)$.

The resulting dimensional wave equation is

$$
\frac{\partial^2 \dot{Q}}{\partial x^2} - \frac{1}{b_r} \frac{\partial \dot{Q}}{\partial x} + \frac{3b_r - \dot{b}}{gA} \frac{\partial^2 \dot{Q}}{\partial x \partial t} + \frac{2b_r}{gA} \frac{\partial \dot{U}_r}{\partial x} \frac{\partial \dot{Q}}{\partial t} - \frac{\partial \dot{Q}}{\partial t} \frac{\partial \dot{U}_r}{\partial x} - \frac{\partial \dot{U}_r}{\partial t} \frac{\partial \dot{Q}}{\partial x} = 0
$$

(5)

The first term on the left-hand side of (5) arises from the pressure gradient, the second is from the pressure gradient and topography, the following two are convective accelerations, the next is the acceleration term and the last is the bed stress.

Scaling the Wave Equation

The most important issue in scaling the wave equation (5) is to determine the magnitude of $\partial \dot{Q} / \partial x = \dot{A} \partial \dot{U} / \partial x$. For weak topography and no tidal flats this is customarily done from (3) under the assumption that the scale for $\dot{x}$ is the inverse wave number $l = \lambda(2\pi)$. The result is that the along-channel velocity scale $U$ is $U = U_0 (g h)^{1/2}$. Intuitively, one would expect that topography of a scale $L < l$ would cause $\partial \dot{U} / \partial x$ to increase. However, the magnitude of $\partial \dot{U} / \partial x$ in a channel flow is constrained by that of $\partial \dot{Q} / \partial t$ which is fixed by the tidal frequency (equation (3)). Thus while topography creates turbulence and may enhance overide generation, the velocity $U_L = \partial \dot{U} / \partial x$ associated with topography of scale $L$ at any given tidal frequency must vary with as $L/\lambda$ so that $\partial \dot{U} / \partial x$ is independent of $L/\lambda$. The scaling used in the weak topography limit ($L \gg l$) is thus appropriate for all $L$; i.e., $\partial \dot{U} / \partial x = U/\lambda$. The velocity for the $n^k$ overide will then scale as $\varepsilon^k U$ in the usual manner. A simple modification of the above scaling is needed to take account for the presence of tidal flats. From (3) and the spreading of wave crests over the tidal flats:

$$
c_0 = \left\{ \frac{B_0}{B_{T_0}} H_0 \right\}^{1/2}
$$

(6)

$$
U = \frac{B_{T_0}}{B_0} c_0 \lambda = \left\{ \frac{B_{T_0}}{B_0} H_0 \right\}^{1/2}
$$

where $B_{T_0}$, $B_0$, and $H_0$ are the total width, width, and depth scales. This choice of $c_0$ will also be justified a posteriori.

That $\partial \dot{U} / \partial x$ in a channel flow scales with the integral scale $l$ regardless of the topography is in marked contrast to the situation in wide estuaries. In these systems, $\partial \dot{U} / \partial x$ may be balanced by $\partial \dot{V} / \partial x$ (where $V$ is the velocity in the cross-channel or $y$ direction) and is not constrained by $\partial \dot{Q} / \partial t$ [Jay, 1990]. Velocity gradients in wide estuaries with small-scale topography may become sufficiently large that the customary series representation of the Stokes drift in terms of particle displacements and velocity gradients [Longuet-Higgins, 1969] does not converge; particle drifts become chaotic [Zimmerman, 1978, 1986]. The present analysis cannot be applied to such systems.

The scaling used to nondimensionalize (4) is

$$
\dot{Q} = U B \dot{H} \dot{Q} \quad \dot{U}_R = U_R \dot{U}_R
$$

$$
\dot{b} = B \dot{b} \quad \dot{h} = H \dot{h} \quad \dot{b}_r = B \dot{b}_r
$$

$$
\dot{t}_s = c_0 U T_b \quad \dot{E} = \varepsilon E_0
$$

$$
\dot{x} = \lambda x \quad \dot{t} = \frac{t}{\varepsilon}
$$

(7)

The resulting nondimensional wave equation is

$$
\frac{\partial^2 \dot{Q}}{\partial x^2} - \frac{R \dot{b}_r}{\omega h} \frac{\partial \dot{t}_s}{\partial x} + \frac{1}{b_r} \frac{\partial \dot{Q}}{\partial x} - \frac{2U_R \dot{U}_R}{c_0} \frac{\partial \dot{Q}}{\partial t} + \frac{\partial^2 \dot{Q}}{\partial t^2} = 0
$$

(8)

where $R = c_0 / (U H)$. The nondimensional number $U_R / c_0$ multiplying the two convective acceleration terms can also be written $\varepsilon U / U$, indicating that the convective acceleration is small relative to the remaining terms. The ratio $R / \omega$ is a measure of the relative strength of friction and acceleration in the wave equation. This varies with the frequency of tidal wave in question but is typically $O(1)$ for the major tidal constituents. Thus the present scaling differs from that of LeBlond [1978], where the acceleration terms were taken to be small. Nonetheless, it will be shown that friction often dominates wave propagation in river estuaries because the effects of acceleration on the wave number are balanced by those of channel convergence. A solution to this form of the wave equation should give $Q$ for the basic tidal motion correct to $O(\varepsilon)$.

The next step in the analysis is representation of the bed stress. The harmonic model formulated herein requires that the quadratic stress term $T_s = c_0 \dot{U} / U$ be replaced by some approximate form that avoids the need to determine the time of the zero crossings of $U$. A number of schemes have been used for this purpose, as reviewed by Godin [1991]. Most common is a Fourier development introduced by Proudman [1953]. This approach suffers from the shortcomings that it does not introduce even overides ($M_0$ and the like) and cannot directly represent the river flow. The Tschebyschev polynomial approach of Dronkers [1964] is used here because of the need to represent river flow effects. The stress is given by

$$
\frac{R}{\omega} \frac{\partial \dot{t}_s}{\partial t} = \frac{R}{\omega} \frac{\partial \dot{t}_s}{\partial t} \left[ \frac{U / U}{U / U} \right]
$$

$$
= \frac{R}{\omega} \frac{1}{\lambda h} \frac{\partial \dot{t}_s}{\partial t} \left[ 1 - \frac{\dot{h}}{h} \right] + O(\varepsilon^2)
$$

$$
= \frac{R}{\omega} \frac{b_r}{\lambda h} \frac{\partial \dot{Q}}{\partial x} + \frac{U_R}{\lambda h} \frac{\partial \dot{Q}}{\partial x} + \frac{c_0}{\lambda h} \frac{\partial \dot{Q}}{\partial x}
$$

(9)

$$
R' = \frac{R}{\pi} \left\{ p_1 + \frac{\dot{U}_R}{U_R} + P_3 \left[ 3U_R^2 + \frac{3}{2} U_R^2 \right] \right\}
$$

$$
R'' = \frac{R}{\pi} \left\{ p_1 + P_2 \left[ \frac{\dot{U}_R}{U_R} + \frac{U_R^2}{U_R} \right] + P_3 \left[ U_R^2 + \frac{3}{2} U_R^2 \right] \right\}
$$
where $U_e$ is the amplitude of $U$, and the time-varying area introduces a finite amplitude nonlinearity. The $p_i$ are Tchebyschev coefficients. They are functions of the ratio $U_p/U$ and therefore vary along the channel. For $U_p = 0$, $p_1 = 0$ and $p_2 = \pi$, as in Proudnikov's [1953] linearization. Far upstream where the current never reverses because the river flow is greater than the tidal flow and $U_p/U \geq 1$, $p_1 = p_2 = 0$ and $p_3 = \pi$ so that $U_p/U = U_p U$. The $p_3$ term is responsible for the $O(1)$ interaction of tidal constituents. Each additional transport $Q_n$ for constituent $k$ adds to $R'/\omega$ a term of the form $A_k \omega / \sqrt{2}$, where $A_k$ is the velocity amplitude for constituent $k$. The present treatment is similar to that of Godin and Gutierrez [1986] for $Q_e/q = 0$. They suggest for analytical purposes $U_p/U = 0.5(mU + 1/mU^3)$ with $m = 0.7$ as an approximation to an error-minimizing form that also contains a term in $U^3$.

Approximate solutions

Approximate solutions to the wave equation (8) with the stress term given by (9) can be found by assuming a harmonic form for the tidal transport: $Q = Re [P(x) e^{i\omega t}]$. The resulting ordinary differential equation is, dropping the overbar on $A$,

$$P_{xx} + T(x) P_x + G(x) P = 0$$

$$G(x) = \frac{1}{bh} \left[ 1 - \frac{ir}{\omega} \frac{1}{h} + 2i \frac{U_k}{c} \frac{dA}{dx} \right]$$

$$T(x) = -\frac{1}{br} \frac{db}{dx} + i \frac{U_k}{c} \left[ \frac{3br}{b} - 1 + \frac{ir}{\omega} \frac{1}{bh^2} \right]$$

where a subscript following a comma indicates differentiation. Two solutions to (10) will be developed here to span the range of physical situations from strong to weak topography. The first covers both weak and strong convergence ($m = 1$, the standard solution), while the second is appropriate to the intermediate, frictionally dominated regime ($m = 2$, the critical solution). Both solutions use a method first developed by Green [1837]: changes first of the independent and then the dependent variable. Green distorted along-channel distance $\hat{x}$ by the inviscid wave speed $(gh)^{1/2}$, so that an inviscid, infinitesimal amplitude wave would move equal distances in the distorted coordinate system in equal time, regardless of depth. Because the inviscid wave speed is altered by the presence of tidal flats and friction, a new independent variable $\hat{y}$ is defined here as $\hat{y} = \hat{x}(gh/br h^m)^{1/2}$ (or $\hat{y} = \hat{x}(b/br h^m)^{1/2}$ in nondimensional variables), with $m = 1$ for the standard solution and $m = 2$ for the critical solution. As is shown below, taking $m = 1$ causes the acceleration and convergence terms in (10) to have constant coefficients. This solution gives Green's Law in the inviscid, weak-topography limit. Taking $m = 2$ causes the friction and convergence terms to have nearly constant coefficients. Both of the solutions defined here exist over the entire range of topographic variability, and either could be used alone, but the changes in the relative importance of the acceleration, friction, and topographic terms in the wave equation as the topographic scale varies cause alterations in solution behavior that are best represented using two separate solutions. The transformation of Prandle and Rahman [1980] ($y = (1 - ir)/10^{1/2} x$ in the present notation) incorporates both acceleration and friction effects on wave speed. It is not used here for several reasons. First a factor of $1/h$ represented explicitly in (9) was included by Prandle and Rahman in $R$. Because of the explicit presence of this factor of $1/h$, use of their transformation would not in the present analysis lead to a constant coefficient equation. Second, such a transformation defines a solution that does not vary smoothly from weak to strong topography. It has a singular point at the intermediate point of critical convergence (defined below), and yet this critically convergent case is precisely where friction is most important. Third, and more physically, topography, not friction, is the primary factor altering the wave speed from its inviscid value. Recognition of the role of convergence/divergence in controlling the nature of the solution suggests the use of two transformations defined here.

The change of independent variable causes the wave equation to take the form

$$P_{yy} + T(x) P_y + G_m P = 0 \quad m = 1, 2$$

$$T_m(y_m) = -\frac{1}{2} \left[ \frac{1}{br} \frac{db}{dy} + \frac{1}{dy_m} \right] + \frac{V_k}{c} \frac{h}{h}$$

$$G_m(y_m) = \left[ \frac{bh^2 + m}{br} \right] \left[ \frac{3br}{b} - 1 + \frac{ir}{\omega} \frac{1}{bh^2} \right]$$

$$Z(y_m) = \exp \left[ -\frac{1}{2} T_m(y_m) dy_m \right]$$

where $m = 1$ and 2 refer to the standard and critical solutions and (15) is the transformed wave equation. Note that (11) to (15) are still general in that no assumption has been made concerning the form of the topography, other than that the channel is narrow. Further progress can only be made by specifying $br$, $b$, and $h$. These may be approximated by exponentials:

$$b_r = B_{(w)} b^m_m$$

$$\beta = B_{(s)} a^{1/2}_m$$

$$k = H_{(s)} h^{1/2}_m$$

with the convergence lengths $\alpha$, $\beta$, and $\gamma$ chosen to best approximate the topography of the system in question. Irregular topography may be dealt with by use of a model with multiple segments, each having different topographic characteristics. Substitution of (16) into (14) gives, after some algebra,
The dimensionless, complex wave number 
\[ \beta_m = \left( \frac{b h_m}{b_T} \right)^{1/2} \frac{\omega}{k_m} \]

where \( k_m \), the real part of \( q_m \), corresponds to the usual inviscid wave number and \( r_m \), the imaginary part of \( q_m \), is the damping modulus [Dronkers, 1964]. \( V(y_m) \) and \( Z(y_m) \) are then

\[ V(y_m) = A e^{i q_m y_m} + B e^{-i q_m y_m} \]

\[ Z(y_m) = \left( b_T b h_m \right)^{1/4} e^{-i \delta_m y_m} \]

where \( A \) and \( B \) are coefficients determined from the boundary conditions and \( \delta_m \) is a real \( O(\epsilon) \) correction to \( k_m l \) that is proportional to \( U_R \) and emerges from integration of the \( O(\epsilon) \) part of \( T(y_m) \) in (15). It serves to retard the incident wave and accelerate the reflected wave in proportion to the strength of the river flow.

The complete form of the solutions for \( Q \) and \( \zeta \) is

\[ Q(y_m, t) = \text{Re} \left( (b_T b h_m)^{1/4} \right. \]

\[ \left. \times \left[ A e^{i(\delta_m - \delta_0) y_m} + B e^{-i(\delta_m + \delta_0) y_m} \right] e^{\omega t} \right] \]

\[ = R_\zeta \left[ (A e^{\gamma_m y_m} + B e^{-\gamma_m y_m}) e^{\omega t} \right] \]

\[ \zeta(y_m, t) = R_\zeta \left[ (b_T b h_m)^{1/4} \right. \]

\[ \left. \times \left[ A e^{i(\delta_m - \delta_0) y_m} + B e^{-i(\delta_m + \delta_0) y_m} \right] e^{\omega t} \right] \]

\[ = \text{Re} \left[ (A e^{\gamma_m y_m} + B e^{-\gamma_m y_m}) e^{\omega t} \right] \]

\[ c^* = \left[ \frac{\Delta_m}{2} - i (q_m - \delta_m) \right] l \]

\[ c^- = \left[ \frac{\Delta_m}{2} - i (q_m + \delta_m) \right] l \]

using continuity to determine \( \zeta \) from \( Q \).

The \((b_T b h_m)^{1/4} = \exp(\Delta/2 \gamma_m) \) and \((b_T b h_m)^{-1/4} = \exp(-\Delta/2 \gamma_m) \) dependences of \( Q \) and \( \zeta \) for the standard solution \((m = 1)\) are for \( b = b_T \) the same as seen in Green’s Law, but this direct topographic dependence or topographic funneling is modified in the presence of friction by the damping modulus \( r_m l \) which also depends on the topography. As is discussed below, topographic dependence through \( r_m l \) often dominates over topographic funneling. It is only in the inviscid, weak topography limit that the standard solution reduces to Green’s Law. The critical solution differs from the standard solution not only in the form of the direct topographic dependence of \( Q \) and \( \zeta \) but also in form of \( q_m l \).

The final step in solution of the wave equation (11) is application of boundary conditions to determine the constants \( A \) and \( B \) in (19), which are, respectively, the complex coefficients of the reflected and incident waves for each segment. Each coefficient specifies an amplitude and phase. Division of the system into \( n \) segments complicates application of the boundary conditions because there are then \( 2n \) constants for the \( n \) incident and reflected waves. These \( 2n \) constants must be determined from two boundary conditions and \( 2n - 2 \) matching conditions between segments by a Gaussian elimination procedure first used in this context by Dronkers [1964]. The matching conditions are linear and guarantee the continuity of \( Q \) and \( \zeta \) between segments. This suffices to guarantee conservation of energy as well [Lighthill, 1978]. One boundary condition (amplitude and phase of the wave at the estuary entrance) is applied at \( x = y_m = 0 \), while the other is applied at the upstream end. There are two possibilities at the upstream boundary. Either the wave dies out uniformly at large \( y_m \), or complete reflection occurs at some finite \( y_m \).
and is \(O(\varepsilon)\). The \(O(1)\) part of \(q_m l\) depends only on three \(O(1)\) terms in \(S_m\) (17). To lowest order, \(q_m l\) is given by

\[
q_m l = \left[ h^{-1} - \frac{[\Delta n l]^2}{2} \right]^{1/2} + O(\varepsilon) \tag{22}
\]

The leftmost term in \(q_m l\) comes from the acceleration term in the wave equation (5). For uniform width and depth with \(b_1 = b\) and no friction, (22) gives the usual long-wave results for the dimensional wave number and celerity \(q = k = \omega(gh)^{1/2}\) and \(c = (gh)^{1/2}\) in the undistorted coordinate system. (Henceforth, explicit identification of the standard and critical solutions by subscript will be made only when necessary for clarity. The subscript \(m\) on \(y, q,\) and \(S\) will also be dropped.) In the presence of tidal flats, \(c = (gb/b_1 h)^{1/2}\), justifying (6). Inclusion of bed stress in (5) introduces an imaginary part to \(q l\) that arises from the right-hand term in (22). The imaginary part of \(q l\) is negative for \(R' > 0\). Choosing the root that gives \(k l > 0\) leads to \(c l < 0\). Thus both incident and reflected waves damp in their respective directions of propagation. The middle term of \(q l\) in (22) is a correction to \(q l\) caused by channel convergence/divergence. It always decreases \(k l\) and increases \(c\).

It is evident from (22) that the \(O(1)\) part of \(q l\) is a function of \(h\) and two nondimensional parameters, \(R' / \omega\) and \(\Delta l\). However, the transformations used in determining the standard and critical solutions are such that \(h\) in the \(O(1)\) part of each solution appears only in a term which is not dominant, and because model segments are short, \(h\) is always \(O(1)\). Therefore the \(O(1)\) behavior of the solution is actually governed in both models by the two nondimensional parameters. \(R' / \omega\) and \(\Delta l\) are respectively the nondimensional ratios of the importance of friction and topographic convergence/divergence to acceleration in the wave equation. The model of Prandle and Rahman [1980] is controlled by two analogous nondimensional numbers.

A general understanding of the behavior of \(q l\) and the incident and reflected waves as a function of the two nondimensional parameters can be obtained from Figures 2 to 5. These plots were constructed using the complete form of \(S\) from (17), with \(\varepsilon = 0.1\), \(U_q = U\), \(U_l = 0.1\), \(b_1 = b\), \(q l = 0\), and \(\Delta l = \Delta l_0 = c l < 0\) (convergent geometry). Constant depth eliminates the distinction between the two solutions. This does not constitute to \(O(1)\) a loss of information, because the dependence of the two models on \(\Omega\) is identical; only the definition of \(\Delta l\) changes. Properties for divergent geometries also differ only at \(O(\varepsilon)\). Consider first the real and imaginary parts of \(q l, k l\) and \(R' / \omega\) (Figures 2a and 2b). The scaling of (7), \(R' / \omega = O(1)\) for the major tidal constituents, \(< 1\) for overides and \(O(10)\) for tidal residual wave components and shelf wave-forced oscillations. Note that \(R'\) is nearly invariant with frequency, because \(R'\) for all waves in a system is set primarily by the strongest tidal constituent(s) and river flow, if this is large.

For \(\Delta l\) very small, \(q l \equiv (1 - i R' / \omega)^{1/2} \equiv 2^{1/4} \left(\cos \frac{\pi}{8} + i \sin \frac{\pi}{8}\right)\), taking \(R' / \omega = 1\). From Figure 2a, \(k l = O(1)\) in this asymptote and only weakly variable over the frequency range of the major tidal species. In the strong-friction (or low frequency) asymptote, however, \(k l\) increases (wave speed decreases) rapidly with increasing \(R' / \omega\). Thus long waves are at least weakly dispersive in the presence of friction. The damping modulus \(R l\) is, in contrast, strongly variable at the weak-friction (high frequency) asymptote and continues to increase slowly with \(R' / \omega\) as the high friction asymptote is approached (Figure 2b).

The quadratic dependence of \(q l\) on changes in total width and depth (the \(\Delta l\) term in (22)) suggest that strong topography, if present, should play a very important role in determining \(q l\) and thus \(c\). From (22) it can be seen that with \(U_q = 0\) and \(m = 1\), the acceleration and convergence terms cancel for \(\Delta l = 2\). This I have called critical convergence. (Inclusion of \(O(\varepsilon)\) effects causes critical convergence to be a function of \(U_q\) and to occur for \(\Delta l > 2\) in Figures 2a and 2b.) As \(\Delta l\) increases beyond critical convergence (supercritical convergence), \(\Delta l\) increasingly dominates \(k l\) and \(R l\). Thus both friction and channel convergence are important in determining wave speed.

The effects of friction and convergence on wave speed for the incident and reflected waves are seen in Figures 3a and 3b. (Subscripts \(i\) and \(r\) indicate here and henceforth the incident and reflected waves.) As long as convergence is subcritical (\(\Delta l < 2\), friction controls \(c / \sigma_0\) for both incident and reflected waves. For \(R' / \omega < 1\), \(c / \sigma_0\) is slightly greater than 1 because of the effect of finite amplitude in accelerating the wave, as is discussed below. For \(R' / \omega O(1)\) the frictional retardation of the tidal wave is less than 10%. This is the same order and opposite
The complex impedance function $I$ (the ratio of $\zeta/Q$) is determined by $q_l$, $\delta l$ and $\Delta l$ and reveals important information about wave propagation. From (21), $I_l = ic^-$ and $I_r = ic^+$ at $z = 0$. For weak convergence and friction, $|I|$ is essentially 1 both for the incident wave (Figure 4) and for the reflected wave as well (not shown). Not surprisingly, friction dominates $|I|$ when friction is strong, while $\Delta l$ dominates when convergence or divergence is strong. $|I_l|$ for the incident wave becomes very large ($Q$ is small relative to $\zeta$) as $|\Delta l|$ becomes large. In this respect also, the incident wave alone for the strong convergence/divergence case mimics a standing wave.

The phase of $I$ is the phase difference $\phi$ between $\zeta$ and $Q$ (Figure 5). Only in a deep channel (weak friction) of uniform width will $\phi$ for the incident wave approach zero, the asymptotic value for an inviscid progressive wave. Nearly inviscid incident (Figure 5) and reflected waves (not shown) both have $\phi \equiv 0^\circ$ just below critical convergence and $\phi \equiv 90^\circ$ just above, with a very abrupt change at critical convergence. Increasing friction acts to make the phase change less abrupt, as was also the case with the wave speed (Figure 3). The range of typical values of $\phi$ for major tidal species is about $10^\circ$ to $45^\circ$ for $\Delta l = 0$.

**Fig. 3.** Log of nondimensional wave celerity $c/c_0$ for convergent geometry for the (a) incident and (b) reflected waves as a function of the logs of $R'/W$ and $-\Delta l$, with solution properties as in Figure 2. River flow and channel topography cause the reflected wave to travel faster in the strong-convergence limit. Strong friction slows incident and reflected waves equally.

sign as the effect of finite amplitude. Thus the small-amplitude inviscid wave speed $c_0$ is a good estimate of $c$, the actual wave speed for a tide in a channel of uniform width and depth. Increasing friction ($R'/W > 1$) slows the wave, but the effect is not dramatic. For weak friction the wave speed $c/c_0$ changes drastically at critical convergence. As $R'/W$ increases, changes in $c/c_0$ are spread out over a larger range of $\Delta l$, a typical effect of friction. Because of the influence of $\delta l$, $c/c_0 > c_0/c_0$, but this effect is subtle unless the convergence and/or river flow are strong.

The rapid wave propagation that occurs in the presence of strong convergence (or divergence) can be rationalized as follows. As $|\Delta l|$ for a section becomes very large, the cross section becomes nearly discontinuous. For a truly discontinuous change in sectional area (i.e., at the end of a channel or the intersection of a very large with a very small channel), incident and reflected waves are essentially equal, and wave propagation occurs at an apparently infinite speed; i.e., the wave phase is nearly constant along the channel leading to the junction. As is discussed below, a uniform, exponential change in cross section (even a supercritical one) does not excite a substantial reflected wave. Thus the behavior of the incident wave alone must asymptotically approach that for a standing wave at a discontinuous change of cross section. Effectively instantaneous propagation can occur only if $c/c_0$ becomes large.
It is in the context of the frictionally dominated, critical solution that the uniform width, strong friction model of LeBlond [1978] is relevant. LeBlond’s scaling resulted in the elimination of the acceleration term in (8) and (22) as small. The result is a diffusive wave equation. More generally, any point on the x axes in Figures 2a and 2b with \( R'/\omega > 10 \) fulfills the requirement of LeBlond’s model that friction dominate over acceleration. From (22) and Figure 5, \( S \) is controlled by its imaginary part, \( k \ell = -r_l \), and \( \phi \) is \( \approx 45^\circ \), halfway between the values for inviscid progressive and standing waves. Friction dominates over acceleration in the vicinity of critical convergence because the effects of acceleration on \( q_l \) are canceled by convergence/divergence. In fact, \( \phi \) in most of the lower right quadrant of Figure 5 is between 30° and 50°. The real part of \( S \) is small throughout this quadrant, because either the frequency is low, the depth is small, or the acceleration effect on the wave number is countered by convergence/divergence. In summary, friction commonly dominates tides in river estuaries more because convergence and acceleration cancel one another than because acceleration is inherently weak. Furthermore, waves for which acceleration effects on \( S \) and \( q_l \) clearly cannot be ignored (e.g., high over­ tides) may still, through the interaction of acceleration and convergence, be frictionally dominated and mimic the behavior of low-frequency, strong-friction waves in a channel of uniform width.

The physical and mathematical bases for the sharp changes in wave properties around critical convergence can be seen from the changes in position of \( S \) in the complex plane as \( R'/\omega \) and \( \Delta \ell \) vary. Consider three cases with varying convergence rates (Figures 6a to 6c). Note first that the imaginary part of \( S \) is always negative, so that \( k \ell \) and \( r_l \) have opposite signs. For weak, sub­critical convergence with \( k \ell \approx R'/\omega \approx 1 \), the real and imaginary parts of \( S \) are nearly equal (Figure 6a). This yields a \( \phi \) of \( \approx 225^\circ \) with \( r_l \) small relative to \( k \ell \). The critical solution is used when \( R'/\omega > 1 \) and/or when acceleration and convergence/­vergence nearly cancel. In either case, \( S \) is almost purely imaginary (Figure 6b). At critical convergence,

\[
q_l = \frac{2R'}{ai, y/,} \cdot \left( \frac{2l}{20} \right)^{1/2}
\]

Because the real part of \( q_l \) is small, the wave speed becomes large (Figure 3). By the same token, a real value of \( \xi \) leads to a largely imaginary value of \( Q \). For \( R'/\omega = 0 \), the real part of \( q_l = ir_l \), and the disturbance ceases to propagate. That is, a supercritically convergent, inviscid wave is impossible, as was noted by Lighthill [1978]. This emphasizes again the importance of friction to any realistic analysis.

### O(\( \epsilon \)) Properties

The \( O(\epsilon) \) properties of \( S \) stem from the convective accelerations and nonlinear, finite-amplitude friction terms. The \( O(\epsilon) \) parts of \( S \) and \( \Delta \ell \) are from (15) and (17) proportional to \( U_f/c = \epsilon U_f/U \) and are influenced by five other nondimensional parameters: \( \beta, b_0, \lambda, y/, \), and \( R'/\omega \). Because \( \Delta \ell \) is real, it modifies only \( k \ell \) and, as was discussed previously, serves to slow the incident wave and accelerate the reflected wave. Because \( q_l \) varies as \( \epsilon^{1/2} \), the \( O(\epsilon) \) terms in \( S \) contribute to both \( k \ell \) and \( r_l \).

There is one \( O(\epsilon) \) feature of the circulation that the perturbation solution presented above does not include: the increase in wave celerity associated with finite amplitude. The classical result [Lamb, 1932, article 187] for inviscid, small-amplitude waves in a uniform channel without tidal flats is that \( c/c_0 = 1 + 3 \epsilon/2 \), or about a 15% increase in \( c \) for \( \epsilon = 0.1 \). In the absence of overdamping topographic effects, neglect of this finite-amplitude correction could lead to an error in calculating high water (even without override effects) of about an hour at a distance of \( \lambda \) from the ocean. This would more than compensate for retardation caused by friction, assuming \( R'/\omega = 1 \). An approximate dispersion relation that accounts for this finite-
amplitude effect on $c$ can be derived as follows. The $O(\varepsilon)$ effect in question arises from nonlinearities in the convective accelerations, pressure gradient, and bed stress. Starting from a dimensional wave equation including these effects, a waveform is then assumed such that $Q \equiv \exp(i(\xi \tilde{x} - \delta \tilde{t}))$. A quasilinearization is used to deal with the nonlinear terms. Some algebra then leads to a modified form of the $O(1)$ part of $S$ for the standard solution

$$S = \left(1 - K_1 \varepsilon\right) \left(\frac{\Delta t}{2} \right)^2 \frac{iR'(1 - \varepsilon)}{\omega h}$$

where $q_0$ is $q$ without finite amplitude correction as calculated from (22). Finite amplitude thus decreases slightly the contribution of friction and acceleration to $S$. It is usually reasonable to approximate $K_1 \approx 3b/b_1$, because the $\gamma + \alpha\gamma q_0$ part of the correction to the acceleration term is only important relative to $3b/b_1$ when acceleration is itself much smaller than the convergence/divergence term. For a channel of uniform width and depth, the inviscid wave speed is $c/c_0 = 1 + (3\varepsilon/2)/(b/b_1)$, which reduces to Lamb's [1932] result in the absence of tidal flats. No corrections are needed to the $O(\varepsilon)$ terms of $S$ (17), because these corrections would be $O(\varepsilon)^2$.

**Along-Channel Variability of the Wave Number**

The validity of the solution presented above depends on, among other factors, the validity of the assumption that the wave number is nearly constant spatially. For weak friction, the topographic convergence rate $\Delta t$ alone determines the wave number, and it is immaterial to $O(1)$ what combination of $\alpha b, \beta b$, and $\gamma b$ makes up $\Delta t$. The presence of friction, however, destroys this symmetry between depth and width changes, because $h$ appears with different powers in the acceleration and friction terms, even after the wave equation has been transformed by a change of coordinate system. Thus several different types of geometry must be considered in evaluating the variability of the wave number. $S$ (or its square root $q_1$) is independent of $y$ to $O(1)$ for $h$ constant, except for long-channel changes in $R'/\omega$. Another important case occurs when $b$ and $b_1$b are constant and equal, so that all topographic variability is caused by changes in $h$. Figures 7a and 7b show the normalized along-channel variability of $q_1$ for the standard and critical solutions, with the realm of use of each indicated, as in Figure 6b. Taking a normalized $q_1$ variability less than 1 as adequate, then only flows that are both strongly frictional and supercritically convergent/divergent are inadequately represented by the solutions derived above. Perhaps the worst case occurs when the change in width is of opposite sign to the change of depth ($b_1 = b, B = -m\gamma$) so that $\Delta t \approx 0$. Even in this circumstance, along-channel variability of $q_1$ is small, except in the extreme of very strong friction and convergence. Variability in $R'/\omega$ occurs in most geometries because of along-channel changes in the tidal velocity and in river flow $U_R$, altering both $R'$ and the $O(\varepsilon)$ parts of $S$. This variability is an important practical consideration only when $U_R$ is greater than $U$ in the upper reaches of tidal rivers.

**A Generalized Form of Green's Law**

The objective of this section is to determine for all values of $R'/\omega$ and $\Delta t$ the dependence $Q$ and $\zeta$ on $b_r$, $h$, and $x$ for both frictional and standard solutions, that is, to generalize Green's Law. In this regard, we are interested only in the amplitudes $|Q|$ and $|\zeta|$ and not phase information contained in the full forms of $Q$ and $\zeta$. Determination of the full topographic dependence of $|Q|$ and $|\zeta|$ requires inclusion of the topographic dependence of $R'$ as well as the effects of topographic funneling, because friction causes $|Q|$ and $|\zeta|$ to depend on the damping modulus $rl$. Algebraic expressions for the normalized topographic dependence of $Q$ may be derived from (21):

$$\frac{b_r d|Q|}{d|Q|} = \left(1 + \frac{2 + rl}{\beta l} \right)$$

$$\frac{h d|Q|}{d|Q|} = \left(4 + \frac{m + rl}{\gamma l} \right)$$

$$\frac{1}{|Q|} \frac{d|Q|}{dx} = \left(2 + \frac{\Delta t}{2 + rl} \right)$$

for $b_r = b$. The expressions in (26) for $b_r$ and $h$ dependences of $Q$ give respectively 1/4 and 1/2 for the standard solution ($m = 1$) and $rl = 0$. Deviations from these values indicate deviations from Green's Law.
Properties for ζ differ from those for Q in (26) in that minus signs are found in front of the first of the two terms in every case. Thus for example, the normalized br dependence of ζ is \((-1/2 + r_1/βl\)). Figures 8 to 13 show normalized dependences of Q and ζ for the incident wave in the standard solution \((m = 1)\) as a function of \(R'/\alpha\) and \(Δl\) for convergent geometry. These plots and (26) may be viewed as a generalization of Green's Law. The plots were constructed with \(br = b\) and \(U_R = 0\). The \(b_r\), \(h\), and \(x\) dependences of the critical solution are identical to those shown in Figures 8 to 13 except as per (26) for the strong-friction and convergence dependence on \(h\), where a factor of \(1/2\) appears instead of \(1/4\). Properties of the reflected wave are identical to those for the incident wave in the weak-friction, weak-topography asymptote and may be understood from (26) and symmetry considerations discussed below in the strong-topography and strong-friction limits. Divergent geometry changes the signs of η, γ and Δ in (26). Thus for example, the \(b_r\) dependence of Q for convergent geometry is the same as that of Q/ for divergent geometry. Plots for the critical solution, divergent geometry, and reflected wave are therefore unnecessary.

It is useful to begin interpretation of Figures 8 to 13 by noting two asymptotic forms of \(r_l\). Consider first the case of weak convergence. As defined in Figure 6a, weak convergence implies that both convergence and friction are small relative to acceleration, and the standard solution is used in the model. From (17), \(r_l = -0.5R'/\alpha(1 - 0.25Δl^2)\) or \(r_l = -0.5R'/\alpha\) for very weak convergence. Topographic funneling dominates the topographic response, and only a slight modification to Green's Law is seen, as in the lower left corners of Figures 8 to 11 where the \(b_r\) variation of Q and ζ approaches \(±1/2\) and the \(h\) variation approaches \(±1/4\); this is also the case for reflected waves.

The strong-topography asymptotic form of \(r_l\) also causes the normalized topographic variations of Q and ζ to have a simple form, because from (24) \(r_l = -1/Δl/2\). Topographic funneling and damping are of the same magnitude. They may therefore either cancel or reinforce one another, and the normalized variation of \(|Q|\) and \(|ζ|\) for \(br = b\) and \(U_R = 0\) accordingly approaches 0 or \(±1\) (top left corner of Figures 8 to 11). There is also, in the absence of river flow, a symmetry between incident/reflected waves in strongly divergent/convergent channels for Q and ζ. For example, ζ/ in a convergent geometry and ζ/ in a divergent geometry with the same wave amplitude and value of \(Δl\) will.
exhibit the same behavior, because both waves encounter the same balance of topography and friction (Figures 12 and 13). The same is true of \( Q \), in divergent geometry in relation to \( Q \), in an equivalent convergent geometry, as summarized in Table 1. This symmetry is, of course, altered by river flow.

The strong friction \( b_f \) and \( h \) variations of \( |Q| \) and \( |\zeta| \) are from (26) dominated by damping rather than topographic funneling for both incident and reflected waves, causing large deviations from Green's Law (lower right-hand corner of Figures 8 to 11). This is true below critical convergence even for \( R' = 0 \) of \( O(10^{-1} \) to 1), which corresponds to the behavior of tide and overtide waves in typical, weakly convergent channels. Because \( R' \) enters incident and reflected waves with opposite sign (equation (26)), the sign of the topographic variations in reflected wave plots is opposite to that in Figures 8 to 13.

The above considerations concerning the behavior of \( \zeta \) and \( Q \) with respect to \( b \), \( h \) and \( x \) apply only to individual incident or reflected waves, not to combinations thereof. Nonetheless, there are numerous circumstances where reflected waves are small and the total wave behavior is determined essentially by the incident wave. These include tides in many rivers, where friction is dominant, and very sharply convergent (in the landward direction) gulfs and bays, where as per (26) friction and topography acting together will sharply damp the reflected wave. Note also in this regard that an incident wave may have any phase difference between \( Q \) and \( \zeta \), even \( 90^\circ \) (Figure 5).

**SUMMARY AND CONCLUSIONS**

Two new solutions to the one-dimensional, tidal long-wave equation for channel flows have been derived that include the effects of friction, smooth topography of arbitrary along-channel scale, finite amplitude, river flow and tidal flats that store water but do not convey momentum. The first, the standard solution, has nearly constant coefficients when the effects of acceleration and topographic convergence are dominant over friction in determining the wave number. This can occur when the topographic convergence is either very strong or very weak. The second, the critical solution, has nearly constant coefficients when friction controls the wave number. The behavior of both solutions is determined (to \( O(1) \)) by two nondimensional parameters. These are the ratio of bed stress to acceleration \( (R'/\omega) \) and the ratio of topographic convergence/divergence rate to acceleration \( (\Delta \omega \alpha) \).

For an inviscid wave, it is immaterial what combination of width and/or depth convergence goes into \( \Delta \omega \). Friction introduces a distinction between width and depth convergence that is, however, small under most circumstances. Both solutions change character at critical convergence, where \( |\Delta \omega \alpha| = 2 \). Inviscid waves do not exist in the strong topography realm above critical convergence \( (|\Delta \omega \alpha| > 2) \). Friction increasingly controls the character of subcritically converging waves as \( R' / \omega \) becomes large. River flow, if present, greatly increases frictional damping.

Topography influences tidal long wave propagation in several ways. First, there is the direct influence of topography on the amplitude of the wave, that is, topographic funneling. For inviscid waves over weak topography (where the topographic scale \( L \) is greater than the inverse wave number \( l \)), this effect is for \( b_f = b \) described by Green's Law: transport \( Q \) varies with width \( b \) and depth \( h \) according to \( Q \equiv b^{1/2} h^{1/4} \) and elevation \( \zeta \) according to \( \zeta \equiv b^{1/2} h^{-1/4} \). It has been shown herein that tidal flats modify topographic funneling so that for weak and strong convergence (the standard solution) \( Q \equiv b_t^{1/4} b_h^{1/4} h^{1/4} \) and \( h \equiv b_t^{1/4} b_h^{1/4} h^{1/4} \), where the total width \( b_t \) differs from the momentum conveying width \( b \) because of the presence of tidal flats. Strong friction (the critical solution) also changes the effect of topographic funneling so that \( Q \) and \( \zeta \) depend on, respectively, \( b_h^{1/2} \) and \( h^{-1/2} \).

Topographic funneling is, however, only part of the picture. The amplitude of a propagating long wave also depends on damping of the wave as described by the imaginary part of the wave number, that is, by the damping modulus \( R' \). Not only is \( R' \) a function of the topography, but if topography and/or friction are strong, the influence of damping on wave amplitude may dominate over or balance that of topographic funneling. In this regard, three regimes may be distinguished. If friction and topog-
Friction, topography, the presence of tidal flats, finite amplitude (as measured by the perturbation parameter \( \epsilon \)), and river flow all influence the real part of the wave number \( k \) \( \pm \delta k \) and the wave propagation speed \( c \). For weak friction and convergence, the effect of friction in slowing the wave is often nearly balanced by the effects of finite amplitude, which scale as \( O(\epsilon) \), are otherwise smaller than the \( O(1) \) effects of friction and topography. River flow modifies wave speed through \( \delta k \) by accelerating the incident wave and deflecting the reflected wave, but this effect is usually weak, because \( U_b \) is normally small relative to \( c \). Tidal flats exercise an influence that may be more important than either finite amplitude or river flow, because tidal flats slow wave propagation by a factor of \( (b_0/b_1)^{1/2} \). This effect arises entirely from mass conservation considerations and is a consequence of spreading of water onto tidal flats as the wave crest advances. For subcritically convergent waves, friction increasingly controls \( k \) as \( R'/\alpha \) becomes large. The wave propagation speed \( c \) of supercritically convergent waves is proportional to \( \Delta \delta \) (\( R'/\alpha \)), which may become very large. Such waves also have a phase difference between \( Q \) and \( \zeta \) that is \( \pi \approx 90^\circ \). Thus a single, incident wave in a strongly convergent topography mimics a standing wave in that the large value of \( c \) causes high water to be nearly simultaneous throughout, and in that high water and slack water are coincident.

The phase difference \( \phi \) between \( Q \) and \( \zeta \) cannot alone be used to infer anything with regard to the presence or absence of a reflected wave. For example, the argument that \( \phi \approx 90^\circ \) implies the existence of a standing wave with equal incident and reflected waves is fallacious for two reasons. First, friction is important in all estuarine systems and dominant in most shallow estuaries. That is, even if perfect reflection from a barrier in an estuary with perfect reflection from a barrier in an estuary with

\[ \frac{\partial^2 u}{\partial y^2} + \frac{\partial^2 u}{\partial z^2} = \frac{1}{\rho} \frac{\partial p}{\partial x} \]

It is effectively a function of \( k \), the damping tends to be independent of amplitude, and the damping is primarily due to friction.

Some mention should also be made of the multiple semidiurnal and diurnal tidal constituents, overtides, and tidal residual phenomena not otherwise considered here. Multiple tidal constituents have not been incorporated into the model at this time, but their inclusion is straightforward, as they interact at \( O(1) \) only through the \( P_3 \) term in (9), so that each constituent damps all others. As the effect on any given constituent is quadratic in the other constituents, only the strongest constituent(s) and the river flow (if strong) are of primary importance in this regard.

Any finite-amplitude long wave will steepen and distort as it propagates, transferring energy at \( O(\epsilon) \) to both zero frequency (the residual flow) and higher harmonics (overtides). Tidal waves of similar frequency will interact at \( O(\epsilon) \) to cause additional overtides and low-frequency motions. Godin [1991] and Parker [1984] have emphasized the role of bed stress nonlinearities in driving both overtides and residuals. There is, however, an important distinction between the effects of strong convergence/divergence on residual versus residual flow generation. The Stokes drift (and therefore the compensating residual elevation and transport) depend for channel flows on \( 0.5|\zeta||\epsilon|\cos \phi \). Like a standing wave with \( \phi \approx 90^\circ \), an incident wave in a strongly divergent/convergent channel generates very little Stokes drift or residual. This is by no means the case with the waves in the same geometry, as these are quadratic in \( \zeta \), but do not depend on \( \phi \). Strong convergence/divergence may be expected to have an important effect on the phase of the residual relative to the basic tidal wave. In this regard, topography may either augment or compete with the effects of tide flats. Speer and Aubrey [1985] have demonstrated that the presence of broad tidal flats changes the phase relationship between the basic tidal wave and its first overtide, such that broad, shallow systems are ebb dominant (overtide reinforces the ebb current) rather than flood dominant, as is typically the case for narrow estuaries. Friedrichs and O. Madsen (manuscript in preparation, 1991) have interpreted long wave behavior in channels with tidal flats in terms of a modified form of \( \epsilon \) (which becomes negative in the presence of broad tidal flats) and a frictional length scale. This behavior may also be understood by considering the additional inhomogeneous forcing terms introduced into the over tide wave equation by the presence of tidal flats. An extension of the present model will consider this problem.

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