Icarus 229 (2014) 11–30

Contents lists available at ScienceDirect

Icarus

journal homepage: www.elsevier.com/locate/icarus

Tidal heating in icy satellite oceans

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ARTICLE INFO

Article history:
Received 14 May 2013
Revised 21 October 2013
Accepted 22 October 2013
Available online 30 October 2013

Keywords:
Satellites, dynamics
Tides, solid body
Jupiter, satellites
Saturn, satellites
Triton

ABSTRACT

Tidal heating plays a significant role in the evolution of many satellites in the outer Solar System; however, it is unclear whether tidal dissipation in a global liquid ocean can represent a significant additional heat source. Tyler (Tyler, R.H. [2008]. Nature 456, 770-772; Tyler, R.H. [2009]. Geophys. Res. Lett. 36, doi:10.1029/2009GL038300) suggested that obliquity tides could drive large-scale flow in the oceans of Europa and Enceladus, leading to significant heating. A critical unknown in this previous work is what the tidal quality factor, Q, of such an ocean should be. The corresponding tidal dissipation spans orders of magnitude depending on the value of Q assumed.

To address this issue we adopt an approach employed in terrestrial ocean modeling, where a significant portion of tidal dissipation arises due to bottom drag, with the drag coefficient O (0.001) being relatively well-established. From numerical solutions to the shallow-water equations including nonlinear bottom drag, we obtain scalings for the equivalent value of Q as a function of this drag coefficient. In addition, we provide new scaling relations appropriate for the inclusion of ocean tidal heating in thermal–orbital evolution models. Our approach is appropriate for situations in which the ocean bottom topography is much smaller than the ocean thickness.

Using these novel scalings, we calculate the ocean contribution to the overall thermal energy budgets for many of the outer Solar System satellites. Although uncertainties such as ocean thickness and satellite obliquity remain, we find that for most satellites it is unlikely that ocean tidal dissipation is important when compared to either radiogenic or solid-body tidal heating. Of known satellites, Triton is the most likely icy satellite to have ocean tidal heating play a role in its present day thermal budget and long-term thermal evolution.

1. Introduction

Tidal heating influences the present-day behavior of some planetary bodies, such as Io (Peale et al., 1979) and Enceladus (Spencer et al., 2006; Howett et al., 2011). It probably also played a role at earlier times elsewhere, including Europa (Hussmann and Spohn, 2004), Ganymede (Showman et al., 1997), Triton (Jankowski et al., 1989), and the Moon (Garrick-Bethell et al., 2010), and may be important in some super-Earth exoplanets (Henning et al., 2009).

For solid bodies, the effects of tides and their associated dissipation are typically calculated assuming a viscoelastic rheology, such as Maxwell, Andrade or Burgers (e.g. Ross and Schubert (1989); Tobie et al. (2005); Efroimsky and Williams (2009); Castillo-Rogez et al. (2011); Nimmo et al. (2012)), though other processes (such as frictional heating, e.g. Nimmo and Gaidos (2002)) may also play a role. For primarily fluid bodies, such as giant planets, a significant component of dissipation is likely to be due to the breaking of internal gravity waves (Ogilvie and Lin, 2004). Lastly, fluid layers on or within solid bodies may also be a source of dissipation. On the Earth it is well-known that tidal dissipation occurs mainly in the oceans (Munk and MacDonald, 1960; Egbert and Ray, 2000; Ray et al., 2001). Global subsurface oceans are thought to occur on at least Europa, Ganymede, Callisto, Titan and perhaps Enceladus (Khurana et al., 1998; Kivelson et al., 2002; Bills and Nimmo, 2011; Iess et al., 2012; Postberg et al., 2011); our focus in this paper is to examine tidal dissipation within such oceans.

In a prescient paper, Ross and Schubert (1989) discussed the possibility of tidal heating on Enceladus arising from turbulent dissipation in a subsurface ocean. More recently, Tyler (2011) expanded an analysis initially developed by Longuet-Higgins (1968) to investigate energy dissipation in tidally-driven satellite oceans. In Tyler (2011), a key free parameter is the linear drag constant α, which can be related to a tidal quality factor Q. It is worth noting that the value of α or Q is *a priori* very poorly known and the total energy dissipation scales linearly with the model's prescribed value.

In this paper, we follow an analysis similar to that of Tyler (2011). However, we depart from his approach in two important respects. First, we provide an estimate for Q using an approach and parameter values developed in studies of the Earth's oceans.

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http://dx.doi.org/10.1016/j.icarus.2013.10.024

Second, we present approximate scaling relationships which allow fluid dissipation to be calculated in a manner analogous to the well-known equations for solid body dissipation (e.g. Segatz et al. (1988); Ross and Schubert (1989); Wisdom (2008)). This will facilitate investigation of long-term satellite evolution, in which the thermal and orbital histories are coupled (e.g. Ojakangas and Stevenson (1989); Hussmann and Spohn (2004); Bland et al. (2009); Meyer et al. (2010); Zhang and Nimmo (2012)).

The rest of this paper is organized as follows. Section 2 reviews the shallow water equations appropriate for flow in global fluid layers on a rotating spherical shell. For clarity, we summarize a semi-analytic solution to these equations, similar to that adopted by Longuet-Higgins (1968) and Tyler (2011), and in addition, explicitly present the method and equations used to calculate quantities such as the average kinetic energy and energy dissipation. Section 3 simplifies this system and carries out an analytical study of the response of a shallow global ocean to tidal forcing, building on the method presented in Section 2. This novel analysis derives approximate scaling relationships for ocean tidal flow and the resulting dissipation under typical icy satellite parameters. The algebra involved can be tedious; to aid clarity, many details have been relegated to Appendices E and F, while Table 4 summarizes the key results. The advantage of these relationships is that they retain the fundamental physical effects while being somewhat simpler to implement than the method presented in Section 2. The results of Section 3 are expressed in terms of an unknown effective (presumably turbulent) viscosity. In Section 4, we present an estimate for this viscosity using a numerical technique based on analogy to frictional ocean dissipation on Earth. We discuss the applications and implications of these results in Section 5. In particular, ocean dissipation is unlikely to be a significant heat source unless the orbital eccentricity is very small; Triton is thus the most likely candidate for a satellite in which ocean tidal dissipation is significant.

2. Ocean tidal dissipation

Here we briefly review the equations of motion for a shallow global satellite ocean. These equations are equivalent to Eqs. (3) and (4) presented in Tyler (2011); we present fully dimensional equations and explicitly expand these equations for the solutions to unknown spherical harmonic coefficients.

The forced, dissipative, shallow water equations on a rotating sphere are (cf. Longuet-Higgins (1968) Eqs. (13.1)–(13.3) or Tyler (2011) Eqs. (3) and (4), noting the sign difference in the tidal potential term)

$$\frac{\partial \vec{u}}{\partial t} + 2\Omega\cos\theta\hat{r}\times\vec{u} = -g\vec{\nabla}\eta - \vec{\nabla}U - \alpha\vec{u} + \nu\nabla^2\vec{u} \tag{1}$$

and

$$\frac{\partial \eta}{\partial t} + h\vec{\nabla}\cdot\vec{u} = 0, \tag{2}$$

where $\vec{u}$ is the radially-averaged, horizontal velocity vector, Ω is the constant rotation rate, $\hat{r}$ is the unit vector in the radial direction, g is the surface gravity, η is the vertical displacement of the surface, and h is the constant ocean depth. The dissipation can be represented as either a linear process with a linear coefficient α or a Navier–Stokes type viscosity with a viscous diffusivity of ν. We assume no radial gradients in our shallow-water model such that the Laplacian operator, ∇^2, has no radial contributions and thus, ν is an effective horizontal diffusivity. U represents the forcing potential due to tides, either eccentricity- or obliquity-related. Eqs. (1) and (2) are valid for incompressible flow under the assumptions that the thickness of the fluid layer is much smaller than the radius of the body ($h \ll R$), the vertical displacement is much smaller than the layer thickness ($\eta \ll h$) and fluid properties are constant (e.g. α and ν). These equations ignore ocean stratification, and thus do not include the effects of internal tides. In addition, they do not include overlying ice shell rigidity, though this effect should be small (Matsuyama, 2012).

2.1. Tidal potentials

For synchronously rotating satellites, such as the regular satellites of Jupiter and Saturn, we are concerned with the ocean flow driven by the eccentricity of the orbit and the obliquity, the tilt of the rotational axis relative to the orbital axis. The forcing tidal potentials can be derived by assuming the planet is a point mass and calculating the resulting gravitational potential at every point on the satellite (cf. Kaula (1964); Murray and Dermott (1999)).

2.1.1. Obliquity tides

The obliquity tidal potential at a point of colatitude θ and longitude ϕ on a synchronously rotating satellite with small obliquity θ_0 (in radians) is a standing wave and can be written as the sum of an eastward and a westward propagating potential (cf. Tyler (2011) Eq. (34))

$$U_{obl} = \frac{-3}{2}\Omega^2R^2\theta_0\sin\theta\cos\theta(\cos(\phi-\Omega t)+\cos(\phi+\Omega t)). \tag{3}$$

We define Laplace spherical harmonics of degree l and order m, Y_l^m, as

$$Y_l^m(\theta,\phi) \equiv \sqrt{\frac{(2l+1)}{4\pi}\frac{(l-m)!}{(l+m)!}}P_l^m(\cos\theta)e^{im\phi} \tag{4}$$

employing a Condon–Shortley phase factor of $(-1)^m$ for $m > 0$. These spherical harmonics are orthogonal under

$$\int_0^{2\pi}\int_0^{\pi} Y_l^m Y_{l'}^{m'*}\ \sin\theta d\theta d\phi = \delta_{l,l'}\delta_{m,m'} \tag{5}$$

where * denotes the complex conjugate and δ is the Kronecker delta.

The obliquity tidal potential thus can be expressed in spherical harmonics as a westward propagating potential $U_{obl,W}$,

$$\begin{aligned} U_{obl,W} &= \frac{3}{2}\sqrt{\frac{2\pi}{15}}\Omega^2R^2\theta_0(e^{i\Omega t}Y_2^1 - e^{-i\Omega t}Y_2^{-1}) \\ &= 2\left(\frac{3}{2}\sqrt{\frac{2\pi}{15}}\Omega^2R^2\theta_0\right)\Re(e^{i\Omega t}Y_2^1) \equiv 2U_{2,W}^1\Re\left(e^{i\Omega t}Y_2^1\right), \end{aligned} \tag{6}$$

and a symmetric eastward propagating potential $U_{obl,E}$ for which the $e^{i\Omega t}$ term is replaced by $e^{-i\Omega t}$ and $U_{2,E}^1 = U_{2,W}^1$.

2.1.2. Eccentricity tides

The eccentricity tidal potential can be expressed as (Kaula, 1964) (cf. Tyler (2011) Eq. (35))

$$\begin{aligned} U_{ecc} = \frac{-3}{4}\Omega^2R^2e[&-(3\cos^2\theta-1)\cos\Omega t + \sin^2\theta(3 \\ &\times\cos 2\phi\cos\Omega t + 4\sin 2\phi\sin\Omega t)]. \end{aligned} \tag{7}$$

For the subsequent analysis, the eccentricity tidal potential can be split into three separate components. There is an axisymmetric component, $U_{ecc,rad}$, (Tyler (2011) calls this the "radial" component)

$$\begin{aligned} U_{ecc,rad} &= \left(3\sqrt{\frac{\pi}{5}}\Omega^2R^2e\cos\Omega t\right)Y_2^0 \\ &= \frac{1}{2}\left(3\sqrt{\frac{\pi}{5}}\Omega^2R^2e\right)(e^{i\Omega t}+e^{-i\Omega t})Y_2^0 \equiv U_2^0(e^{i\Omega t}+e^{-i\Omega t})Y_2^0 \end{aligned} \tag{8}$$

There is also an asymmetric librational component, $U_{ecc,lib}$, that can be split into a westward propagating potential