Tidal asteroseismology: Kepler’s KOI-54

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ABSTRACT

We develop a general framework for interpreting and analysing high-precision light curves from eccentric stellar binaries. Although our methods are general, we focus on the recently discovered Kepler system KOI-54, a face-on binary of two A stars with \( e = 0.83 \) and an orbital period of 42 days. KOI-54 exhibits strong ellipsoidal variability during its periastron passage; its light curve also contains \( \sim 20 \) pulsations at perfect harmonics of the orbital frequency, and another \( \sim 10 \) non-harmonic pulsations. Analysis of such data is a new form of asteroseismology in which oscillation amplitudes and phases rather than frequencies contain information that can be mined to constrain stellar properties. We qualitatively explain the physics of mode excitation and the range of harmonics expected to be observed. To quantitatively model observed pulsation spectra, we develop and apply a linear, tidally forced, non-adiabatic stellar oscillation formalism including the Coriolis force. We produce temporal power spectra for KOI-54 that are semi-quantitatively consistent with the observations. Both stars in the KOI-54 system are expected to be rotating pseudo-synchronously, with resonant non-axisymmetric modes providing a key contribution to the total torque; such resonances present a possible explanation for the two largest-amplitude harmonic pulsations observed in KOI-54, although we find problems with this interpretation. We show in detail that the non-harmonic pulsations observed in KOI-54 can be explained by non-linear three-mode coupling. The methods developed in this paper can be generalized in the future to determine the best-fitting stellar parameters given pulsation data. We also derive an analytic model of KOI-54’s ellipsoidal variability, including both tidal distortion and stellar irradiation, which can be used to model other similar systems.

Key words: asteroseismology – hydrodynamics – waves – binaries: close – stars: oscillations.

1 INTRODUCTION

The recently discovered Kepler system KOI-54 (Welsh et al. 2011, henceforth W11) is a highly eccentric stellar binary with a striking light curve: a 20-h 0.6 per cent brightening occurs with a periodicity of 41.8 days, with lower amplitude perfectly sinusoidal oscillations occurring in between. Such observations were only possible due to the unprecedented photometric precision afforded by Kepler. W11 arrived at the following interpretation of these phenomena: during the periastron passage of the binary, each of its two similar A stars is maximally subjected to both its companion’s tidal force and radiation field. The tidal force causes a prolate ellipsoidal distortion of each star known as the equilibrium tide, so that the resulting perturbations to both the stellar cross-section and the emitted stellar flux produce a change in the observed flux. Along with the effects of irradiation, this then creates the large brightening at periastron. Secondly, the strong tidal force also resonantly excites stellar eigenmodes during periastron, which continue to oscillate throughout the binary’s orbit due to their long damping times; this resonant response is known as the dynamical tide. W11 successfully exploited KOI-54’s periastron flux variations, known traditionally as ellipsoidal variability, by optimizing a detailed model against this component of KOI-54’s light curve (Orosz & Hauschildt 2000). In this way, W11 were able to produce much tighter constraints on stellar and orbital parameters than could be inferred through traditional spectroscopic methods alone. W11 also provided data on the dynamical tide oscillations. 30 such pulsations were reported, of which roughly two-thirds have frequencies at exact harmonics of the orbital frequency. It is the analysis of these and similar future data that forms the basis of our work.

In close binary systems, tides provide a key mechanism to circularize orbits and synchronize stellar rotation with orbital motion. An extensive literature exists on the theory of stellar tides (e.g. Zahn

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1975; Goodman & Dickson 1998; Witte & Savonije 1999). We have synthesized this theoretical formalism, together with other aspects of stellar oscillation theory, in order to model the dynamical tide of KOI-54 as well as to provide a framework for interpreting other similar systems.

The methods we have begun to develop are a new form of asteroseismology, a long-standing subject with broad utility. In traditional asteroseismology, we observe stars in which internal stellar processes (e.g. turbulent convection or the kappa mechanism) drive stellar eigenmodes, allowing them to achieve large amplitudes (Christensen-Dalsgaard 2003). In this scenario, modes ring at their natural frequencies irrespective of the excitation mechanism. The observed frequencies (and linewidths) thus constitute the key information in traditional asteroseismology, and an extensive set of theoretical techniques exist to invert such data in order to infer stellar parameters and probe different aspects of stellar structure (Unno et al. 1989).

In tidal asteroseismology of systems like KOI-54, however, we observe modes excited by a periodic tidal potential from an eccentric orbit; tidal excitation occurs predominantly at \( f = 2 \) (Section 3.2). Since orbital periods are well below a star’s dynamical time-scale, it is \( g \) modes (buoyancy waves) rather than higher frequency \( p \) modes (sound waves) that primarily concern us. Furthermore, since modes in our case are forced oscillators, they do not ring at their natural eigenfrequencies, but instead at pure harmonics of the orbital frequency. (We discuss non-harmonic pulsations in Section 6.5.) It is thus pulsation amplitudes and phases that provide the key data in tidal asteroseismology.

This set of harmonic amplitudes and phases in principle contains a large amount of information. One of the goals for future study is to determine exactly how the amplitudes can be optimally used to constrain stellar properties, e.g. the radial profile of the Brunt–Väisälä frequency. In this work, however, we focus on the more modest tasks of delineating the physical mechanisms at work in eccentric binaries and constructing a coherent theoretical model and corresponding numerical method capable of quantitatively modelling their dynamical tidal pulsations.

This paper is organized as follows. In Section 2 we give essential background on KOI-54. In Section 3 we give various theoretical results that we rely on in later sections, including background on tidal excitation of stellar eigenmodes (Section 3.2), techniques for computing disc-averaged observed flux perturbations (Section 3.3) and background on including the Coriolis force using the traditional approximation (Section 3.4). In Section 4 we use these results to qualitatively explain the pulsation spectra of eccentric stellar binaries, particularly what governs the range of harmonics excited. In Section 5 we confront the rotational evolution of KOI-54’s stars, showing that they are expected to have achieved a state of stochastic pseudo-synchronization.

In Section 6 we present the results of our more detailed modelling. This includes an analytic model of ellipsoidal variability (Section 6.1), the effects of non-adiabaticity (Section 6.2), the effects of fast rotation (Section 6.3) and a preliminary optimization of our non-adiabatic method against KOI-54’s pulsation data (Section 6.4). We show in Section 6.5 that the observed non-harmonic pulsations in KOI-54 are well explained by non-linear three-mode coupling, and perform estimates of instability thresholds, which may limit the amplitudes modes can attain. We also address whether the highest-amplitude observed harmonics in KOI-54 are signatures of resonant synchronization locks in Section 6.6. We present our conclusions and prospects for future work in Section 7.

Table 1. List of KOI-54 system parameters as determined by W11. Selected from table 2 of W11. The top rows contain standard observables from stellar spectroscopy, whereas the bottom rows result from W11′s modelling of photometric and RV data. Symbols either have their conventional definitions or are defined in Section 3.1. Note that W11′s convention is to use the less-massive star as the primary.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
<th>Error</th>
<th>Unit</th>
</tr>
</thead>
<tbody>
<tr>
<td>Observations</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( T_1 )</td>
<td>8500</td>
<td>200</td>
<td>K</td>
</tr>
<tr>
<td>( T_2 )</td>
<td>8800</td>
<td>200</td>
<td>K</td>
</tr>
<tr>
<td>( L_2/L_1 )</td>
<td>1.22</td>
<td>0.04</td>
<td></td>
</tr>
<tr>
<td>( v_{\text{rot1}} \sin i_1 )</td>
<td>7.5</td>
<td>4.5</td>
<td>km s(^{-1})</td>
</tr>
<tr>
<td>( v_{\text{rot2}} \sin i_2 )</td>
<td>7.5</td>
<td>4.5</td>
<td>km s(^{-1})</td>
</tr>
<tr>
<td>([\text{Fe/H}]_1)</td>
<td>0.4</td>
<td>0.2</td>
<td></td>
</tr>
<tr>
<td>([\text{Fe/H}]_2)</td>
<td>0.4</td>
<td>0.2</td>
<td></td>
</tr>
<tr>
<td>Light curve/RV modelling</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( M_2/M_1 )</td>
<td>1.025</td>
<td>0.013</td>
<td></td>
</tr>
<tr>
<td>( P_{\text{orb}} )</td>
<td>41.8051</td>
<td>0.0003</td>
<td>d</td>
</tr>
<tr>
<td>( e )</td>
<td>0.8342</td>
<td>0.0005</td>
<td></td>
</tr>
<tr>
<td>( \omega )</td>
<td>36.22</td>
<td>0.90</td>
<td></td>
</tr>
<tr>
<td>( i )</td>
<td>5.52</td>
<td>0.10</td>
<td></td>
</tr>
<tr>
<td>( a )</td>
<td>0.395</td>
<td>0.008</td>
<td>au</td>
</tr>
<tr>
<td>( M_1 )</td>
<td>2.32</td>
<td>0.10</td>
<td>( M_\odot )</td>
</tr>
<tr>
<td>( M_2 )</td>
<td>2.38</td>
<td>0.12</td>
<td>( M_\odot )</td>
</tr>
<tr>
<td>( R_1 )</td>
<td>2.19</td>
<td>0.03</td>
<td>( R_\odot )</td>
</tr>
<tr>
<td>( R_2 )</td>
<td>2.33</td>
<td>0.03</td>
<td>( R_\odot )</td>
</tr>
</tbody>
</table>

A few weeks prior to the completion of this manuscript, we became aware of a complementary study of KOI-54’s pulsations (Fuller & Lai 2011).

2 BACKGROUND ON KOI-54

Table 1 gives various parameters for KOI-54 resulting from W11’s observations and modelling efforts. Table 2 gives a list of the pulsations W11 reported, including both frequencies and amplitudes.

2.1 Initial rotation

KOI-54’s two components are inferred to be A stars. Isolated A stars are observed to rotate much more rapidly than e.g. the Sun, with typical surface velocities of \( \sim 100 \text{ km s}^{-1} \) and rotation periods of \( \sim 1 \text{ d} \) (Adelman 2004). This results from their lack of a significant convective envelope, which means they experience less-significant magnetic braking, allowing them to retain more of their initial angular momentum as they evolve on to the main sequence. We thus operate under the assumption that both component stars of KOI-54 were born with rotation periods of roughly \( P_{\text{birth}} \approx 1 \text{ d} \).

2.2 Rotational inclination

W11 constrained both stars’ rotation, via line broadening, to be \( v_{\text{rot}} \sin i = 7.5 \pm 4.5 \text{ km s}^{-1} \) (the same for both stars). Using the mean values of \( R_1 \) and \( R_2 \) obtained from W11’s modelling, we can
translate this into the following constraints on rotation periods (in days):

\[ 9.2 < P_1 / \sin i_1 < 37 \quad \text{and} \quad 9.8 < P_2 / \sin i_2 < 39, \]

where \((i_1, i_2)\) and \((P_1, P_2)\) are the rotational angular momentum inclinations with respect to the observer and the stellar rotation periods, respectively. If we assume that tidal interactions cause both stellar rotation periods to be approximately equal to the pseudo-synchronous period of \(P_{ps} \sim 1.8 \text{d} \) derived in Section 5, we can constrain \(i_1\) and \(i_2\):

\[ 2.8 < i_1 < 11^\circ \quad \text{and} \quad 2.6 < i_2 < 11^\circ. \]

W11 obtained \(i_{\text{orb}} = 5\text{.}52 \pm 0.10\) by fitting the light curve’s ellipsoidal variation together with radial velocity (RV) measurements, so the constraints just derived are consistent with alignment of rotational and orbital angular momenta,

\[ i = i_{\text{orb}} = i_1 = i_2. \]  

(1)

Tides act to drive these three inclinations to be mutually parallel or antiparallel, so such an alignment once achieved is expected to persist indefinitely. In order to simplify the analytical formalism as well as reduce the computational expense of modelling the observed pulsations, we will adopt equation (1) as an assumption for the rest of our analysis.

### 3 THEORETICAL BACKGROUND

In this section we review various heterogeneous theoretical results that we rely on in later sections. In Section 3.1 we summarize the conventions and definitions used in our analysis. In Section 3.2 we review the theory of tidally forced adiabatic stellar eigenmodes. Later (Section 4), we use this formalism to explain qualitative features of the light curves of eccentric binaries like KOI-54. We also use adiabatic normal modes to compute tidal torques (Section 5 and Appendix C), as well as to perform a non-linear saturation calculation (Section 6.5). However, our detailed quantitative modelling of the observations of KOI-54 utilizes a non-adiabatic tidally forced stellar oscillation method that we introduce and employ in Section 6.

In Section 3.3 we summarize how perturbed quantities at the stellar photosphere, specifically the radial displacement and Lagrangian flux perturbation, can be averaged over the stellar disc and translated into an observed flux variation. Lastly, in Section 3.4 we review the traditional approximation, a way of simplifying the stellar oscillation equations in the presence of rapid rotation.

#### 3.1 Conventions and definitions

We label the two stars as per Table 1, consistent with W11; note that the primary/star 1 is taken to be the smaller and less massive star. In the following, we focus our analysis on star 1, since the results for star 2 are similar. We assume that both stars’ rotational angular momentum vectors are perpendicular to the orbital plane (Section 2.2), and work in spherical coordinates \((r, \theta, \phi)\) centred on star 1 where \(\theta = 0\) aligns with the system’s orbital angular momentum and \(\phi = 0\) points from star 1 to star 2 at periastron.

We write the stellar separation as \(D(t)\) and the true anomaly as \(f(t)\), so that the position of star 2 is \(\mathbf{D} = (D, \pi/2, f)\). We write the semimajor axis as \(a\) and the eccentricity as \(e\). The angular position of the observer in these coordinates is \(\hat{n}_o = (\theta_o, \phi_o)\), where these angles are related to the traditional inclination \(i\) and argument of periastron \(\omega\) by (Arras et al. 2011)

\[ \theta_o = i \quad \text{and} \quad \phi_o = \frac{\pi}{2} - \omega \mod 2\pi. \]  

(2)

The orbital period (angular frequency) is \(P_{\text{orb}}(\Omega_{\text{orb}})\), while a rotation period (angular frequency) is \(P_i(\Omega_i)\). The effective orbital frequency at periastron is

\[ \Omega_{\text{peri}} = \frac{d f}{d t|_{f=0}} = \frac{\Omega_{\text{orb}}}{1 - e} \sqrt{\frac{1 + e}{1 - e}}, \]  

(3)

which is \(\Omega_{\text{peri}} = 20 \times \Omega_{\text{orb}}\) for KOI-54. The stellar dynamical frequency is

\[ \omega_{\text{dyn}} = \sqrt{\frac{G M}{R^3}}, \]  

(4)

which is \(\omega_{\text{dyn}} \approx 1.1 \text{ rad h}^{-1}\) for KOI-54’s stars.

#### 3.2 Tidal excitation of stellar eigenmodes

Although we ultimately use an inhomogeneous, non-adiabatic code including the Coriolis force to model the pulsations in eccentric binaries (Section 6.2), the well-known normal mode formalism provides an excellent qualitative explanation for many of the features in the light-curve power spectra of systems such as KOI-54. Here we will review the salient results of this standard theory; we demonstrate their application to KOI-54 and related systems in Sections 4 and 5. The remainder of the paper after Section 5 primarily uses our non-adiabatic method described in Section 6.2.
Working exclusively to linear order and operating in the coordinates specified in the previous section, we can represent the response of star 1 (and similarly for star 2) – all oscillation variables such as the radial displacement \(\xi\), the Lagrangian pressure perturbation \(\Delta p\), etc. – to a perturbing tidal potential by a spatial expansion in normal modes and a temporal expansion in orbital harmonics (e.g. Kumar, Ao & Quataert 1995):

\[
\delta X = \sum_{nlm} A_{nlm}(\xi X_n(r) e^{-im\phi}) Y_{lm}(\theta, \phi).
\]

(5)

Here, \(2 \leq l \leq \infty\) and \(-l \leq m \leq l\) are the spherical harmonic quantum numbers, and index the angular expansion; \(|n| < \infty\) is an eigenfunction’s number of radial nodes, and indexes the radial expansion;\(^1\) and \(|k| < \infty\) is the orbital Fourier harmonic number, which indexes the temporal expansion.

Each \((n, l, m)\) pair formally corresponds to a distinct mode, although the eigenspectrum is degenerate in \(n\) for now we are ignoring the influence of rotation on the eigenmodes. Each mode has associated with it a set of eigenfunctions for the various perturbation variables, e.g. \(\xi, \Delta p, \) etc., as well as a frequency \(\omega_{nl}\) and a damping rate \(\gamma_{nl}\). For stars and modes of interest, \(\gamma_{nl}\) is set by radiative diffusion; see the discussion after equation (17). Fig. 1 gives a propagation diagram for a stellar model consistent with W11’s mean parameters for Table 1 (Table 1). The frequencies of g modes behave asymptotically for \(n \gg 0\) and hence \(\omega_{nl} \ll \omega_{dyn}\) as (Christensen-Dalsgaard 2003)

\[
\omega_{nl} \sim \frac{\omega_{nl}}{n},
\]

(6)

where \(\omega_{nl} \approx 4\text{ rad h}^{-1}\) for KOI-54’s stars.

The amplitudes \(A_{nlm}\) appearing in equation (5) each represent the pairing of a stellar eigenmode with an orbital harmonic. Their values are set by the tidal potential, and can be expressed analytically as:

\[
A_{nlm} = \frac{2 \ell ! Q_{nl} \tilde{X}_{lm}^{1} W_{lm} \Delta_{nlm}^{1}}{E_{nl}}.
\]

(7)

The coefficients appearing in equation (7) are as follows.

(i) The tidal parameter \(\ell !\) is given by

\[
\ell ! = \frac{M_{2}}{M_{1}} \left(\frac{R_{1}}{D_{peri}}\right)^{l+1},
\]

(8)

where \(D_{peri} = a(l - e)\) is the binary separation at periastron. This factor represents the overall strength of the tide; due to its dependence on \(R_{1}/D_{peri}\), which is a small number in cases of interest, it is often acceptable to consider only \(l = 2\).

(ii) The linear overlap integral \(Q_{nl}\) (Press & Teukolsky 1977), given by

\[
Q_{nl} = \frac{1}{M_{1}R_{1}^{2}} \int_{0}^{R_{1}} l \left(\xi_{nl}(r) + (l + 1)\xi_{nl}(r)\right) \rho r^{l+1} dr
\]

\[
= \frac{1}{M_{1}R_{1}^{2}} \int_{0}^{R_{1}} \delta_{nl} r^{l+2} dr
\]

\[
= -\frac{R_{1}}{GM_{1}} 2l + 1 \frac{1}{4\pi} \delta_{nl} (R_{1}),
\]

(9)

represents the spatial coupling of the tidal potential to a given eigenmode; it is largest for modes with low \(|n|\) and hence for eigenfrequencies close to the dynamical frequency \(\omega_{dyn} = \sqrt{GM_{1}/R_{1}^{3}}\), but falls off as a power law for \(|n| \gg 0\).

Conventionally, \(n > 0\) corresponds to p modes while \(n < 0\) corresponds to g modes; however, since we are mostly concerned with g modes in this paper, we will report g mode \(n\) values as >0.

(iii) We define our mode normalization/energy \(E_{nl}\) as

\[
E_{nl} = 2\left(\frac{\omega_{nl}^{2}R_{1}}{GM_{1}}\right) \int_{0}^{R_{1}} \left(\xi_{nl}^{2} + (l + 1)\xi_{nl}^{2}\right) \rho r^{l+1} dr,
\]

(10)

where \(\xi_{h}\) is the horizontal displacement (Christensen-Dalsgaard 2003).

(iv) The unit-normalized Hansen coefficients \(\tilde{X}_{lm}^{1}\) are the Fourier series expansion of the orbital motion (Murray & Dermott 1999), and are defined implicitly by

\[
\left(\frac{D_{peri}}{D(r)}\right)^{l+1} e^{-im\phi} = \sum_{k=-\infty}^{\infty} \tilde{X}_{lm}^{1}(e) e^{-ik\Omega_{obs}}.
\]

(11)

They are related to the traditional Hansen coefficients \(X_{lm}^{1}\) by \(X_{lm}^{1} = \tilde{X}_{lm}^{1}/(1 - e)^{l+1}\) and satisfy the sum rule:

\[
\sum_{k=-\infty}^{\infty} \tilde{X}_{lm}^{1}(e) = 1,
\]

(12)

which can be verified using equation (11). (An explicit expression for \(X_{lm}^{1}\) is given in equation A6.) The Hansen coefficients represent the temporal coupling of the tidal potential to a given orbital harmonic. They peak near \(k_{\text{peak}} \sim m\Omega_{\text{peri}}/\Omega_{\text{orb}}\) but fall off exponentially for larger \(|k|\).

(v) The Lorentzian factor \(\Delta_{nlm}\) is

\[
\Delta_{nlm} = \frac{\omega_{nl}^{2}}{\left(\omega_{nl}^{2} - \sigma_{lm}^{2}\right) - 2i\gamma_{nl}/\Omega_{\text{obs}}},
\]

(13)

where \(\sigma_{lm} = k\Omega_{\text{obs}} - m\Omega_{i}\) and represents the temporal coupling of a given harmonic to a given mode. When its corresponding mode/harmonic pair approach resonance, i.e. \(\omega_{nl} \approx \sigma_{lm}\), \(\Delta_{nlm}\) can become very large; its maximum, for a perfect resonance, is half the simple harmonic oscillator quality factor, \(q_{nl} = \Delta_{nl} = \omega_{nl}/2\gamma_{nl}\).

(vi) \(W_{lm}\) is defined in equation (A3) and represents the angular coupling of the tidal potential to a given mode; it is non-zero only for \(mod(l + m, 2) = 0\). In particular, \(W_{2,1} = 0\), meaning that \(l = 2, |m| = 1\) modes are not excited by the tidal potential.

(vii) To calculate the quasi-adiabatic damping rate \(\gamma_{nl}\) within the adiabatic normal mode formalism,\(^2\) we average the product of the thermal diffusivity \(\chi\) with a mode’s squared wavenumber \(k^{2}\), weighted by the mode energy:

\[
\gamma_{nl} = \frac{\int_{0}^{R_{1}} k^{2} \left[\xi_{nl}^{2} + (l + 1)\xi_{nl}^{2}\right] \rho r^{l+1} dr}{\int_{0}^{R_{1}} \left[\xi_{nl}^{2} + (l + 1)\xi_{nl}^{2}\right] \rho r^{l+1} dr},
\]

(14)

where the thermal diffusivity \(\chi\) is

\[
\chi = \frac{16\pi}{3k\rho^{3/2}} c_{g}.
\]

The cut-off radius \(r_{c}\) is determined by the minimum of the mode’s outer turning point and the point where \(\omega_{n}\) is \(2\pi\) (Christensen-Dalsgaard 2003), where the thermal time is

\[
t_{\text{therm}} = \frac{pc_{g} T}{g F}.
\]

(16)

When this cut-off is restricted by the mode period intersecting the thermal time, so that strong non-adiabatic effects are present inside the mode’s propagation cavity, the mode becomes a travelling wave at the surface, and the standing wave/adiabatic normal mode

\(^2\) We only use this approximate method of calculating damping rates when employing the adiabatic normal mode formalism; our non-adiabatic method introduced in Section 6.2 fully includes radiative diffusion.

\(^1\)Conventionally, \(n > 0\) corresponds to p modes while \(n < 0\) corresponds to g modes; however, since we are mostly concerned with g modes in this paper, we will report g mode \(n\) values as >0.

\(^3\)The result we obtain is not well approximated by a series expansion, and the time scale we interpret the cut-off radius as \(t_{\text{therm}}\) in the case of the tidal asteroseismology of Kepler-11 radial modes.
approximation becomes less realistic. This begins to occur at a frequency (in the rest frame of the star) of \(50 \times \Omega_{\text{orb}}\) for KOI-54, as can be seen in Fig. 1. Fortunately, our calculations involving the normal mode formalism (Sections 5 and 6.5) centre primarily on low-order modes that are firmly within the standing wave limit.

The g-mode damping rate scales roughly as

\[
\gamma_n \sim \gamma_0 n^4 \sim \gamma_0 \left( \frac{\omega_n}{\omega_{\text{damp}}} \right)^4,
\]

where we have used the asymptotic g-mode frequency scaling from equation (6). In the standing wave regime, i.e. for \(\omega_{\text{damp}} \lesssim 50 \times \Omega_{\text{orb}}\), we find that \(\gamma_0 \sim 1 \text{ Myr}^{-1}\) and \(s \sim 4\). This large value for \(s\) results from the fact that most of the damping occurs at the surface, and the cut-off radius is limited by the outer turning point where the mode frequency intersects the Lamb frequency. As the mode frequency declines, the cut-off radius moves outwards towards smaller Lamb frequency and stronger damping, as can be seen in Fig. 1. Without this behaviour of the turning point, we would expect \(s \sim 2\) since \(k^2 \propto n^2\) in equation (14).

### 3.3 Observed flux perturbation

Throughout this work, perturbations to the emitted flux \(\Delta F\) are understood to be bolometric, i.e. integrated over the entire electromagnetic spectrum. We correct for Kepler’s bandpass to first order as follows. We define the bandpass correction coefficient \(\beta(T)\) as the ratio of the bandpass-corrected flux perturbation \((\Delta F/F)_{\text{hpc}}\) to the bolometric perturbation \((\Delta F/F)\), so that

\[
\frac{(\Delta F)}{F}_{\text{hpc}} = \beta(T) \left( \frac{\Delta F}{F} \right).
\]

We assume Kepler is perfectly sensitive to the wavelength band \((\lambda_1, \lambda_2) = (400, 865)\) nm (Koch et al. 2010), and is completely insensitive to all other wavelengths. Then \(\beta(T)\) is given to first order by

\[
\beta(T) \approx \int \frac{\delta B_\lambda}{\delta \ln T} \, d\lambda, \quad 4 \int \frac{\delta B_\lambda}{\lambda} \, d\lambda,
\]

where \(B_\lambda(T)\) is the Planck function. Using W11’s mean parameters for KOI-54 (Table 1), we have \(\beta(T_1) = 0.81\) and \(\beta(T_2) = 0.79\). Note that employing \(\beta\) alone amounts to ignoring bandpass corrections due to limb darkening. We have also ignored the fact that in realistic atmospheres, the perturbed specific intensity depends on perturbations to gravity in addition to temperature; this is a small effect, however, as shown e.g. in Robinson, Kepler & Nather (1982).

For completeness, we transcribe several results from Pfahl et al. (2008), which allow a radial displacement field \(\xi_i\) and a Lagrangian perturbation to the emitted flux \(\Delta F\), both evaluated at the stellar surface, to be translated into a corresponding disc-averaged observed flux perturbation \(\delta J\), as seen e.g. by a telescope (Dziembowski 1977). While an emitted flux perturbation alters the observed flux directly, a radial displacement field contributes by perturbing a star’s cross-section.\(^3\)

Given \(\xi_i\) and \(\Delta F\) expanded in spherical harmonics as

\[
\xi_i = \sum_{l=0}^{\infty} \sum_{m=-l}^{l} \xi_{i,lm}(t) Y_{lm}(\theta, \phi),
\]

\[
\Delta F = \sum_{l=0}^{\infty} \sum_{m=-l}^{l} \Delta F_{lm}(t) Y_{lm}(\theta, \phi),
\]

we can translate these into a fractional observed flux variation \(\delta J/J\) to first order by

\[
\frac{\delta J}{J} = \sum_{l=0}^{\infty} \sum_{m=-l}^{l} \left[ 2(b_l - c_l) + \beta(T) b_l \frac{\Delta F_{lm}(t)}{F(R)} \right] Y_{lm}(\theta_0, \phi_0),
\]

where the disc-integral factors are

\[
b_l = \int_0^1 \mu P_l(\mu) h(\mu) \, d\mu,
\]

\[
c_l = \int_0^1 \left[ 2\mu^2 \frac{dP_l}{d\mu} - (\mu - \mu^2) \frac{d^2 P_l}{d\mu^2} \right] h(\mu) \, d\mu,
\]

\(P_l(\mu)\) is a Legendre polynomial, and \(h(\mu)\) is the limb-darkening function, normalized as \(\int_0^1 \mu h(\mu) \, d\mu = 1\). For simplicity, we use Eddington limb darkening for all of our analysis, with \(h(\mu) = 1 + 3\mu/2\); \(b_l\) and \(c_l\) in this case are given in Table 3 for \(0 \leq l \leq 5\).

Since \(Y_{2,\pm2} \propto \sin^2 \theta\) and \(Y_{2,0} \propto (3\cos^2 \theta - 1)\), and since KOI-54 has \(\theta_0 = i \approx 5.5\) (Table 1), equation (22) shows that \(m = \pm 2\) eigenmodes are a factor of \(\sim 200\) less observable than \(m = 0\) modes. It is thus likely that nearly all of the observed pulsations in KOI-54 have \(m = 0\); the exceptions may be F1 and F2, as we discuss in Section 6.6.

3. A horizontal displacement field \(\xi_x\) produces no net effect to first order – its influence cancels against perturbations to limb darkening, all of which is included in equation (22).
3.4 Rotation in the traditional approximation

Stellar rotation manifests itself in a star’s corotating frame as the fictitious centrifugal and Coriolis forces (Unno et al. 1989). The centrifugal force directly affects the equilibrium structure of a star, which can then consequently affect stellar oscillations. Its importance, however, is characterized by $(\Omega/\omega_{\text{dyn}})^2$, which is $\sim 10^{-2}$ for rotation periods and stellar parameters of interest here (Section 2). As such we neglect rotational modification of the equilibrium stellar structure (Ipser & Lindblom 1990).

The Coriolis force, on the other hand, affects stellar oscillations directly. Given a frequency of oscillation $\sigma$, the influence of the Coriolis force is characterized by the dimensionless rotation parameter $q$ given by

$$q = 2\Omega / \sigma,$$  

where large values of $|q|$ imply that rotation is an important effect that must be accounted for. Note that for simplicity we assume rigid-body rotation throughout. For the pulsations observed in KOI-54’s light curve (Table 2), assuming both stars rotate at near the pseudo-synchronous rotation period $P_{\text{orb}} \sim 1.8$ d discussed in Section 5.1 and that $m = 0$ (justified in the previous section), $q$ ranges from 0.5 for $k = 10$ to 1.5 for $k = 30$. Thus, lower harmonics fall in the non-perturbative rotation regime, where rotation is a critical effect that must be fully included.

The ‘traditional approximation’ (Chapman & Lindzen 1970) greatly simplifies the required analysis when the Coriolis force is included in the momentum equation. In the case of $g$ modes, it is applicable for

$$l \gtrsim \frac{2 R}{q H_p} \left( \frac{\Omega}{|N|} \right)^2,$$  

where $H_p = \rho g / \rho_p$ is the pressure scale height; outside of the convective cores of models we are concerned with in this work (where $g$ modes are evanescent anyway), equation (26) is well satisfied whenever rotation is significant. Here we will give a brief overview of the traditional approximation; we refer to Bildsten, Ushominsky & Cutler (1996) for a more thorough discussion.

The traditional approximation changes the angular Laplacian, which occurs when deriving the non-rotating stellar oscillation equations, into the Laplace tidal operator $L^\mu_{\nu}$. (Without the traditional approximation, the oscillation equations for a rotating star are generally not separable.) It is thus necessary to perform the partial expansion of oscillation variables in eigenfunctions of $L^\mu_{\nu}$, known as the Hough functions $H^\mu_{\nu}(\mu)$, rather than associated Legendre functions; the azimuthal expansion is still in $e^{i\mu\theta}$. The eigenvalues of $L^\mu_{\nu}$ are denoted $\lambda$, and depend on $\mu$, the azimuthal wavenumber. In the limit that $q \to 0$, the Hough functions become ordinary (appropriately normalized) associated Legendre functions, while $\lambda \to (l+1)$.

We present the inhomogeneous, tidally driven stellar oscillation equations in the traditional approximation in Appendix A2. The principal difference relative to the standard stellar oscillation equations is that terms involving $l(l+1)$ either are approximated to zero or have $l(l+1) \to \lambda$. This replacement changes the effective angular wavenumber, e.g. since the primary $\lambda$ for $m = 0$ increases with increasing rotation, fast rotation leads to increased damping of $m = 0$ g modes at fixed frequency, as discussed in Section 6.3.

For strong rotation, $|q| > 1$, the Hough eigenvalues $\lambda$ can be both positive and negative. The case of $\lambda > 0$ produces rotationally modified traditional $g$ modes, which evanesce for $\cos^2 \theta > 1/|q|^2$. (Rossby waves or $r$ modes are also confined near the equator and have a small positive value of $\lambda$.) Instead, for $\lambda < 0$, polar modes are produced that propagate near the poles for $\cos^2 \theta > 1/|q|^2$, but evanesce radially from the surface since they have an imaginary Lamb frequency $S_{\lambda} = \lambda^{1/2} c/r$ (as explained further in Fig. 1).

4 QUALITATIVE DISCUSSION OF TIDAL ASTEROSEISMOLOGY

It is helpful conceptually to divide the tidal response of a star into two components: the equilibrium tide and the dynamical tide (Zahn 1975). Note that in this section we will again use the normal mode formalism described in Section 3.2, even though our subsequent more detailed modelling of KOI-54 uses the inhomogeneous, non-adiabatic formalism introduced in Section 6.2.

4.1 Equilibrium tide

The equilibrium tide is the ‘static’ response of a star to a perturbing tidal potential, i.e. the large-scale perturbation due to differential gravity from a companion. In terms of light curves, the equilibrium tide corresponds to ellipsoidal variability (along with the irradiation component of this effect discussed in Appendix B1). In the case of an eccentric binary, this manifests itself as a large variation in the observed flux from the binary during periastron. KOI-54’s equilibrium tide was successfully modelled in W11, enabling precise constraints to be placed on various stellar and orbital parameters (Table 1).

In terms of the normal mode formalism developed in Section 3.2, the equilibrium tide corresponds to the amplitudes from equation (7) tied to large overlaps $Q_{ln}$ and large Hansen coefficients $X_{lnk}^\mu$; in other words, to pairings of low-|$l$| modes with low-|$k$| orbital harmonics. The Lorentzian factor $D_{\text{ad}}$ is typically $\sim 1$ for the equilibrium tide since it is not a resonant phenomenon.

In practice, however, it is much simpler and more convenient to use other mathematical formalisms to model the equilibrium tide, like taking the zero-frequency stellar response as in Appendix B2, or filling Roche potentials as in W11’s simulations. We show in Section 6.1 that our simple analytical treatment of the equilibrium tide verifies the results from the sophisticated simulation code employed in W11.

4.2 Dynamical tide

The dynamical tide, on the other hand, corresponds to resonantly excited pulsations with frequencies equal to harmonics of the orbital frequency, $k \Omega_{\text{orb}}$. W11 observed at least 21 such harmonics (Table 2).

For a circular orbit, the tidal potential has all its power in the $k = \pm 2$ orbital harmonics; in this case the only modes that can be resonantly excited are those with frequencies close to twice the Doppler-shifted orbital frequency: $\omega_{\text{ad}} \approx 2|\Omega_{\text{orb}} - \Omega|$; this is typically only a single mode. This corresponds to the fact that the Hansen coefficients from equation (11) become a Kronecker delta

Table 3. First six disc-integral factors $b_l$ and $c_l$ from equations (23) and (24) for linear Eddington limb darkening, $h(\mu) = 1 + (3/2)\mu$. 

<table>
<thead>
<tr>
<th>$l$</th>
<th>$b_l$</th>
<th>$c_l$</th>
<th>$l$</th>
<th>$b_l$</th>
<th>$c_l$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>1</td>
<td>0</td>
<td>3</td>
<td>1/16</td>
<td>3/4</td>
</tr>
<tr>
<td>1</td>
<td>17/24</td>
<td>17/12</td>
<td>4</td>
<td>-1/48</td>
<td>-5/12</td>
</tr>
<tr>
<td>2</td>
<td>13/40</td>
<td>39/20</td>
<td>5</td>
<td>-1/128</td>
<td>-15/64</td>
</tr>
</tbody>
</table>

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at zero eccentricity: $\tilde{X}_{lm}^k(0) = \delta_{ll}^k$. However, for a highly eccentric orbit, the distribution of power in the Hansen coefficients, and hence the stellar response, can be much broader; as a result a wide array of different harmonics can be excited, allowing for a rich pulsation spectrum.

Mode excitation due to a tidal harmonic $k\Omega_{\text{orb}}$ is modulated by the Doppler-shifted frequency $\sigma_{lm} = k\Omega_{\text{orb}} - m\Omega_\star$. However, the frequencies at which modes are observed to oscillate, viewed from an inertial frame, are indeed pure harmonics of the orbital frequency, $k\Omega_{\text{orb}}$.\footnote{Welsh et al. (2011) incorrectly attributed non-harmonic pulsations to rotational splitting; we return to non-harmonic pulsations in Section 6.5.} We demonstrate this mathematically in Appendix A; intuitively, although a driving frequency experiences a Doppler shift upon switching to a star’s corotating frame, the star’s response is then Doppler shifted back upon observation from an inertial frame. In general, any time a linear system is driven at a particular frequency, it then also oscillates at that frequency, with its internal structure reflected only in the oscillation’s amplitude and phase.

Whether a given mode is excited to a large amplitude is contingent on several conditions – essentially all the terms in equation (7). First, the overall strength of the tide, and hence the magnitude of observed flux variations, is determined by the tidal factor $\varepsilon_1$ from equation (8). The dominant multipole order is $l = 2$, so we have $\varepsilon_2 = (M_2/M_1)(R_*/D_{\text{peri}})^3$, where $D_{\text{peri}} = a(1 - e)$ is the binary separation at periastron, and we are focusing our analysis on star 1. For KOI-54, $\varepsilon_2 \approx 4 \times 10^{-3}$ for both stars.

Next, the strength of a mode’s resonant temporal coupling to the tidal potential is given by the Lorentzian factor $\Delta_{\text{res}}$, in equation (13). Since this factor is set by how close a mode’s frequency is to the nearest orbital harmonic, its effect is intrinsically random. The degree of resonance has an enormous effect on a mode’s contribution to the observed flux perturbation, meaning that modelling the dynamical tide amounts on some level to adjusting stellar and system parameters in order to align eigenfrequencies against orbital harmonics so that the array of Lorentzian factors conspires to reproduce observational data.

Moreover, given a single observed pulsation amplitude together with theoretical knowledge of the likely responsible mode, i.e. the first four factors in equation (7), equating theoretical and observed pulsation amplitudes in principle yields direct determination of the mode’s eigenfrequency, independently of the degree of resonance. This line of reasoning of course neglects the considerable theoretical uncertainties present, but serves to illustrate tidal asteroseismology’s potential to constrain stellar parameters.

Despite the inherent unpredictability, a light curve’s Fourier spectrum is still subject to restrictions imposed primarily by the remaining two factors in equation (7). These terms, the linear overlap integral $Q_{nl}$ and the unit-normalized Hansen coefficient $\tilde{X}_{lm}^k(\varepsilon)$ (respectively equations 9 and 11), restrict the range in $k$ over which harmonics can be excited; Fig. 2 shows profiles of both. As discussed in Section 3.2, $Q_{nl}$ peaks for modes with frequencies near the dynamical frequency of the star $\omega_{\text{dyn}}$, and falls off as a power law in frequency, whereas $\tilde{X}_{lm}^k$ peaks for harmonics near $m\Omega_{\text{peri}}/\Omega_{\text{orb}}$ and falls off for higher $k$:

$$Q_{nl} \propto \omega_{nl}^c, \quad \omega_{nl} \ll \omega_{\text{dyn}},$$

$$\tilde{X}_{lm}^k(\varepsilon) \propto \exp(-k/r) \quad |k| \gg m\Omega_{\text{peri}}/\Omega_{\text{orb}}.$$  \hspace{1cm} (27)

The power-law index $p$ is 11/6 for $g$ modes in stars with a convective core and a radiative envelope or vice versa (Zahn 1970), and for KOI-54’s eccentricity and $l = 2$, we find $r \sim 15$.

As a result, modes that can be excited are those with frequencies in the intervening region between the peaks of $Q_{nl}$ and $X_{lm}^k$, i.e. $m|\Omega_{\text{peri}}| < \omega_{nl} < \omega_{\text{dyn}}$. This is a necessary but not sufficient condition; Fig. 3 shows the product $Q_{nl}X_{lm}^k(\varepsilon)$ at various eccentricities with stellar parameters as well as the periastron distance $D_{\text{peri}}$ fixed to the mean values

$$\tilde{X}_{lm}^k(\varepsilon) \propto \exp(-k/r) \quad |k| \gg m\Omega_{\text{peri}}/\Omega_{\text{orb}}.$$  \hspace{1cm} (28)

\footnote{Welsh et al. (2011) incorrectly attributed non-harmonic pulsations to rotational splitting; we return to non-harmonic pulsations in Section 6.5.}
in W11, and hence with fixed tidal parameter $e_I$ (but consequently allowing the orbital period to vary). Although a chance close resonance can yield a large Lorentzian factor $A_{\text{link}}$, excitation of modes far from the peak of $Q_{\text{lin}} \chi_{lm}^2$, becomes less and less likely, since this quantity falls off sharply, especially towards larger $|k|$. 

There are two other constraints on the range of harmonic pulsations that can be excited. First, the eigenmode density for g modes scales asymptotically as

$$\frac{dn}{dk} \sim \frac{1}{k^2 \Omega} \gamma,$$

which shows that fewer modes exist at higher $k$. This can be seen by the spacing of points (which denote normal modes) in Figs 2 and 3, as well as by the spacing of peaks in Fig. 6. This further limits the number of harmonics that can be excited at large $k$, in addition to the exponential decay of the Hansen coefficients discussed earlier, and thus effectively shifts the curves in Fig. 3 towards lower $k$.

In addition, the Lorentzian factor $A_{\text{link}}$ is attenuated by mode damping $\gamma_{nl}$, which is set by radiative diffusion for high-order g modes. Damping becomes larger with decreasing $g$-mode frequency due to increasing wavenumber; an asymptotic scaling is given in equation (17). Because the Lorentzian response is proportional to $\gamma_{nl}^2$ at perfect resonance, the amplitudes of lower frequency modes/harmonics are diminished by increased damping, in addition to the power-law decay of the tidal overlap. This effect is critical for understanding the influence of rotation on light-curve power spectra, as we investigate in Section 6.3.

### 4.3 Pulsation phases

Pulsation phases in eccentric binaries are essential information which should be fully modelled, in addition to the pulsation amplitudes reported in W11. For simplicity, we focus on a particular harmonic amplitude $A_{\text{link}}$ from equations (5) and (7) and assume it results from a close resonance so that $\omega_{nl} \approx \sigma_{km} = k \Omega_{\text{orb}} - m \Omega_s$, assuming without loss of generality that $\sigma_{km} > 0$. We can then evaluate its phase $\psi_{\text{link}}$ relative to periastron, modulo $\pi$ (since we are temporarily ignoring the real part of the amplitude, which could introduce a minus sign), as

$$\psi_{\text{link}} = \text{arg}(A_{\text{link}}) = \frac{\pi}{2} - \arctan\left( \frac{\delta \omega/\gamma_{nl}}{2 \gamma_{nl} \sigma_{km}} \right) \mod \pi$$

where $\delta \omega = \omega_{nl} - \sigma_{km}$ is the detuning frequency.

For a near-perfect resonance, where $\delta \omega \ll \gamma_{nl}$, $\psi_{\text{link}}$ approaches $\pi/2$ (modulo $\pi$). However, if eigenmode damping rates are much smaller than the orbital frequency, then this intrinsic phase cannot be real. This is the case for K01-54, where $\Omega_{\text{orb}}/\gamma_{nl} > 10^3$ for modes of interest. Indeed, theoretically modelling the largest-amplitude 90th and 91st harmonics of K01-54 assuming they are $m = 0$ modes requires only $|\delta \omega/\gamma_{nl}| \approx 20$, so that even these phases should be within $\pm 1$ per cent of zero (modulo $\pi$).

The phase of the corresponding observed harmonic flux perturbation can be obtained from equation (31) by further including the phase of the spherical harmonic factor in the disc-averaging formula, equation (22):

$$\text{arg}(\delta \Omega_{nl}/J) = \psi_{\text{link}} + m \phi_0.$$  

Summing over the complex conjugate pair, the observed time dependence is then $\cos[k \Omega_{\text{orb}} t - (\psi_{\text{link}} + m \phi_0)]$, where $t \approx \text{Co}r$ corresponds to periastron. However, the sign of $k$ in this formalism is unknown; equivalently, whether the pulsation is prograde or retrograde (Appendix C1) cannot be determined in this way. Thus, if the observed pulsation’s (cosine) phase is $\delta$, the comparison to make is

$$\delta = \pm (\psi_{\text{link}} + m \phi_0) \mod \pi.$$

Nonetheless, since we have argued that $\psi_{\text{link}} \approx 0$, this becomes

$$\delta \approx \pm m \phi_0 \mod \pi.$$  

Given determination of $\phi_0$ (related to the argument of periastron $\omega$ by equation 2) by modelling of RV data or ellipsoidal variation, the phase of a resonant harmonic thus directly gives the mode’s value of $|m|$ (which is very likely 0 or 2 for tidally excited modes, since $l = 2$ dominates). For K01-54, phase information on harmonics 90 and 91 would thus determine whether they result from resonance locks, as discussed in the next section, or are simply chance resonances. Furthermore, knowing $|m|$ allows $m \phi_0$ to be removed from equation (33), yielding the pulsation’s damping-to-detuning ratio.

However, the preceding treatment is only valid if the eigenfunction itself has a small phase: although eigenfunctions are purely real for adiabatic normal modes, local phases are introduced in a fully non-adiabatic calculation, as in Section 6.2. Thus, equations (32)–(34) are only applicable in the standing wave regime, where the imaginary part of the flux perturbation is small relative to the real part. In the travelling wave regime, the local wave phase near the surface becomes significant, and can overwhelm the contribution from global damping; see Section 6.2. For K01-54, this corresponds to $|k|$ below $\sim 30$, although this depends on the rotation rate (Section 6.3).

### 5 ROTATIONAL SYNCHRONIZATION IN K01-54

Here we will discuss a priori theoretical expectations for K01-54’s stars’ rotation. Later, in Section 6.3, we will compare the results derived here with constraints imposed by the observed pulsation spectrum.

#### 5.1 Pseudo-synchronization

In binary systems, the influence of tides causes each component of the binary to eventually synchronize its rotational and orbital motions, just as with Earth’s moon. Tides also circularize orbits, sending $e \rightarrow 0$, but the circulation time-scale $\tau_{\text{circ}}$ is much greater than the synchronization time-scale $\tau_{\text{sync}}$; their ratio is roughly given by the ratio of orbital to rotational angular momenta:

$$\frac{\tau_{\text{sync}}}{\tau_{\text{circ}}} \sim \frac{L_s}{L_s / \gamma} = \frac{\mu a^2 \Omega_{\text{orb}} / \Omega_s}{\mu a^2 \Omega_{\text{orb}} / \Omega_s} \sqrt{1 - e^2}$$

$$\sim \frac{a}{R} \left( \frac{M_s R_e^2}{I_s} \right) (1 - e)^3,$$

where $I_s$ is the stellar moment of inertia, $\mu$ is the reduced mass, and we have assumed for simplicity that the stars rotate at the periastron frequency $\Omega_{\text{peri}}$ (Section 3.1). For K01-54, this ratio is $\sim 10^7$.

Due to the disparity of these time-scales, a star in an eccentric binary will first synchronize to a pseudo-synchronous period $\tau_{\text{ps}}$, defined as a rotation period such that no average tidal torque is exerted on either star over a sufficiently long time-scale. If only the torque due to the equilibrium tide is used, and thus eigenmode resonances are neglected, then only one unique pseudo-synchronous
period exists, $P_{rot}^s$, as derived in Hut (1981) and employed in W11. Its value for KOI-54 is (equation C13)

$$P_{rot}^s = (2.53 \pm 0.01) \text{d}.$$  

Inclusion of eigenmode resonances, however, complicates the situation. Fig. 4 shows the secular tidal torque (averaged over one rotation period) for star 1 of KOI-54 plotted as a function of rotation frequency/period including contributions from both the equilibrium and dynamical tide. Although the general torque profile tends to zero at $P_{rot}^s$, numerous other roots exist (displayed as vertical lines), where the torque due to a single resonantly excited eigenmode of the dynamical tide cancels against that due to the equilibrium tide. To produce this plot, we directly evaluated the secular tidal torque (Appendix C1) using an expansion over the quadrupolar adiabatic normal modes of a MESA stellar model (Paxton et al. 2011) with parameters set by W11’s mean values for star 1 of KOI-54 (Table 1). In our calculation we include both radiative (Section 3.2) and turbulent convective damping (Willems, Deloye & Kalogera 2010), but neglect rotational modification of the eigenmodes.

Next, of the many zeroes of the secular torque available, which are applicable? Continuing with the assumption that KOI-54’s stars were born with rotation periods of $P_{birth} \sim 1 \text{d}$ (Section 2.1), with the same orientation as the orbital motion, one might naively posit that the first zero encountered by each star should constitute a pseudo-synchronous period – it is an ostensibly stable spin state since small changes to either the stellar eigenmodes (via stellar evolution) or the orbital parameters (via circularization and orbital decay) induce a restoring torque. This is the basic idea behind a resonance lock (Witte & Savonije 1999).

However, this conception of resonance locking neglects two important factors. First, although the dynamical and equilibrium tidal torques may cancel, their energy deposition rates do not (in general); see Appendix C1. Thus, during a resonance lock the orbital frequency must continue to evolve, allowing other modes to come into resonance, potentially capable of breaking the lock. Secondly, as shown by Fuller & Lai (2011), it is necessary that the orbital frequency not evolve so quickly that the restoring torque mentioned earlier be insufficient to maintain the resonance lock. This restricts the range of modes capable of resonantly locking, introducing an upper bound on their inertial-frame frequencies and hence their orbital harmonic numbers (values of $k$ in our notation).

Consequently, pseudo-synchronyization is in reality a complicated and dynamical process, consisting of a chain of resonance locks persisting until eventually $e \to 0$ and $P_s = P_{orb}$. Such resonance lock chains were studied in much greater detail by Witte & Savonije (1999) for eccentric binaries broadly similar to KOI-54. As a result of the inherent complexity, a full simulation of KOI-54’s orbital and rotational evolution is required in order to address the phenomenon of resonance locking and to derive theoretical predictions for the stars’ spins. To perform such simulations, we again expanded the secular tidal torque and energy deposition rate over normal modes (detailed in Appendix C1) using two MESA stellar models consistent with W11’s mean parameters for KOI-54’s two stars. We then numerically integrated the orbital evolution equations (Witte & Savonije 1999) assuming rigid-body rotation. We did not include the Coriolis force, nor did we address whether the eigenmode amplitudes required to produce the various resonance locks that arise are stable to non-linear processes (Section 6.5).

Our simulations indicate that both stars should have reached pseudo-synchronization states with rotation periods of $P_{rot}^s \sim 1.8 \text{d}$; we discuss the synchronization time-scale in more detail in Section 5.2. These periods are $\sim 30$ per cent faster than Hut’s value of $P_{rot}^s = 2.5 \text{d}$. The pseudo-synchronization mechanism that operates is stochastic in nature, in which the dynamical tide’s prograde resonance locks balance the equilibrium tide in a temporally averaged sense. This result appears independently of the initial rotation rates used; in other words, it is an attractor.

As described above, when a star is locked in resonance, it is the torque from a single highly resonant eigenmode that acts to oppose the equilibrium tide’s non-zero torque. Such a high-amplitude mode should be easily observable. At first glance, this line of reasoning seems to provide a natural explanation for the presence of the large-amplitude 90th and 91st observed harmonics in KOI-54 (F1 and F2 from Table 2), namely that each is the photometric signature of the highly resonant eigenmode that produces a resonance lock for its...
Figure 5. Simple analytic model of ellipsoidal variability detailed in Appendix B, including both the equilibrium tide (red dotted lines) and ‘reflection/irradiation (blue dashed lines) components of the light curve. We used the best-fitting parameters from W11 in all three panels (Table 1, except that we show two examples of edge-on orientations (ignoring the possibility of eclipses) in (b) and (c), effectively presenting KOI-54’s light curve as it would be observed from different angles. [We used \( \omega = (80^\circ, 20^\circ) \) for (b, c) in order to demonstrate the asymmetric light curves possible depending on the binary’s orientation.] Panel (a) reproduces W11’s modelling and the data for KOI-54 to \(~\)20 per cent. Our analytic model is easily applicable to many other systems. In Section 6.3 we show that the dynamical tidal response, ignored here, may be larger than that due to the equilibrium tide for edge-on systems.

6 RESULTS

6.1 Ellipsoidal variation

Fig. 5(a) shows our simple model of KOI-54’s ellipsoidal variation; we adopted the best-fitting parameters from W11’s modelling (Table 1) to produce our light curve. Our irradiation (Appendix B1, blue dashed line) and equilibrium tide (Appendix B2, red dotted line) models are larger than W11’s results by 24 and 14 per cent, respectively. The shapes of both curves are, however, essentially indistinguishable from W11’s much more detailed calculations.

We attribute the small difference between our results and those of W11 to our simple model of the bandpass correction (equation 19) which ignores bandpass variations due to limb darkening. Such details could easily be incorporated into our analytical formalism, however, by introducing a wavelength-dependent limb-darkening function \( h_{\lambda}(\mu) \) in the disc integrals in equations (23) and (24) (Robinson et al. 1982). We thus believe that the models provided in Appendix B should be quite useful for modelling other systems like KOI-54, due in particular to their analytic simplicity.

We also show in Figs 5(b) and (c) what KOI-54’s equilibrium tide and irradiation would look like for two edge-on orientations, demonstrating the more complicated, asymmetric light-curve morphologies possible in eccentric binaries (see also the earlier work by Kumar et al. 1995). Future searches for eccentric binaries using Kepler and other telescopes with high-precision photometry should allow for the wide range of light-curve shapes shown in Fig. 5. We note, however, that the dynamical tidal response, ignored in this section, may be larger than that due to the equilibrium tide for edge-on systems, as we show in Section 6.3.

6.2 Non-adiabatic inhomogeneous method

Thus far our theoretical results have primarily utilized the tidally forced adiabatic normal mode formalism. Although this framework provides excellent intuition for the key physics in eccentric binaries, it is insufficient for producing detailed theoretical light curves, since this necessitates tracking a star’s tidal response all the way to the

\[ t_{\text{sync}} \sim I \int \frac{d\Omega}{\tau(\Omega)} \times \Omega. \]
photosphere where non-adiabatic effects are critical. To account for this, we employ the non-adiabatic inhomogeneous formalism originally used by Pfahl et al. (2008) (Appendix A1), which we have extended to account for rotation in the traditional approximation (Appendix A2).

Rather than decompose the response of the star into normal modes, the inhomogeneous method directly solves for the full linear response of the star to an external tidal force produced by a companion at a given forcing frequency. Given a stellar model, an orbital period, a set of orbital harmonics to act as driving frequencies and a rigid-body rotation period, we solve the numerical schemes described in Appendix A2 for each star. This determines the various physical perturbation variables of the star as a function of radius, such as the radial displacement and the flux perturbation. For stars of interest we can safely ignore perturbations to the convective flux, so the only non-adiabatic effect is that produced by radiative diffusion.

Fig. 6 shows the surface radial displacement $\xi_r$ and Lagrangian emitted flux perturbation $\Delta F$ computed on a fine frequency grid, temporarily ignoring rotation; normal mode frequencies correspond to the resonant peaks in these curves. The surface radial displacement should approach its equilibrium tide value as the driving frequency tends to zero. Quantitatively, we find that this is true for orbital harmonics $k \lessapprox 30$; note that in the units employed in Fig. 6, this equilibrium tide value for $\xi_r$ is $\xi_r/R_{\text{phot}} = 1$ (Appendix B2).

The surface flux perturbation shown in Fig. 6, on the other hand, more clearly demonstrates the three qualitatively different regimes possible at the surface. First, the weakly damped standing wave regime, $k \gtrapprox 30$, is characterized by strong eigenmode resonances and all perturbation variables having small imaginary parts. In Fig. 1, this corresponds to the outer turning point, where the mode frequency intersects the Lamb frequency, lying inside the point where the mode frequency becomes comparable to the thermal frequency, so that the mode becomes evanescent before it becomes strongly non-adiabatic.

Next, the travelling wave regime, $5 \lessapprox k \lessapprox 30$, arises when modes instead propagate beyond where the mode and thermal frequencies become comparable, leading to rapid radiative diffusion near the surface. In the travelling wave limit, resonances become severely attenuated as waves are increasingly unable to reflect at the surface, and all perturbation variables have comparable real and imaginary parts (not including their equilibrium tide values).

Lastly, just as with the radial displacement, the flux perturbation also asymptotes to its overdamped equilibrium tide/von Zeipel value of $(\Delta F/F)_{\text{eq}} = -(l+2)(\xi_r/R_{\text{phot}})^{2}$ (Appendix B2) in the low-frequency limit, which is $|\Delta F/F|_{\text{eq}} = |1 - 2l| = 4$ in the units employed in Fig. 6. Quantitatively, however, this only occurs for $k \lessapprox 5$. At first glance, this suggests that the equilibrium tide modelling of KOI-54 in W11 and Fig. 5 is invalid, since the equilibrium tide in KOI-54 has orbital power out to at least $k \sim 30$ (as can be seen e.g. in the plot of the Hansen coefficients for KOI-54’s eccentricity in Fig. 2).

Fortunately, as we describe in the next section, including rotation with a face-on inclination effectively stretches the graph in Fig. 6 towards higher $k$, e.g. for $P_\ast = 2.0$ d, we find the equilibrium tide/von Zeipel approximation to hold for $k \lessapprox 30$, justifying the simplifications used in W11 and Appendix B2, although this may not apply for edge-on systems.

### 6.3 Effect of rotation on the dynamical tidal response

The most important effects of rotation in the context of tidal asteroseismology can be seen in Fig. 7. Here we show the predicted flux perturbation for KOI-54 as a function of orbital harmonic $k$ for four different rotation periods, having subtracted the equilibrium tide (Appendix B2) to focus on resonant effects.\(^5\) In Fig. 7(a) we use KOI-54’s face-on inclination of $i = 5.5$, while in Fig. 7(b) we use an inclination of $i = 90^\circ$ to illustrate how a system like KOI-54 would appear if seen edge-on; all other parameters are fixed to those from W11’s modelling (and are thus not intended to quantitatively reproduce the data; see Fig. 8 for an optimized model). The details of which specific higher harmonics have the most power vary as rotation changes mode eigenfrequencies, moving eigenmodes into and out of resonance. Nonetheless, several qualitative features can be observed.

For KOI-54’s actual face-on orientation, as in Fig. 7(a), rotation tends to suppress power in lower harmonics. This can be understood as follows. Primarily $m = 0$ modes are observable face-on (Section 3.3). At fixed driving frequency $\sigma$, as the stellar rotation frequency $\Omega_\ast$, and hence the Coriolis parameter $q = 2\Omega_\ast\sigma$, increases in magnitude, $m = 0$ g modes become progressively confined to the stellar equator (Section 3.4). As a result, these rotationally modified modes angularly couple more weakly to the tidal potential, diminishing their intrinsic amplitudes. Moreover, equatorial compression

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\(^5\) We assume that rotation is in the same sense as orbital motion throughout this section.
Figure 7. Influence of rotation on the light-curve temporal power spectra of eccentric binaries. The eight leftmost plots show theoretical power spectra using fiducial stellar models and orbital parameters consistent with W11’s mean values (Table 1), each as a function of orbital harmonic $k$. The top row uses KOI-54’s known face-on inclination of $i = 5.5^\circ$; the bottom row instead uses an edge-on inclination of $i = 90^\circ$, showing how a system like KOI-54 would appear if viewed edge-on. The vertical axis has different scales in the two rows. The effects of rotation on the stellar pulsations are included using the traditional approximation (Section 3.4). The equilibrium tide (Appendix B2) has been subtracted in order to focus on resonant effects. We have not included negative-$\lambda$ modes, as discussed in the text. (Note that parameters here have not been optimized to reproduce KOI-54’s light curve; see Fig. 8 for such a model.) The four leftmost columns show four different rigid-body rotation periods. (a) For a perfectly face-on orientation, only $m = 0$ modes can be observed (Section 3.3); however, larger rotation rates lead to equatorial compression of lower frequency $m = 0$ g modes, which increases the effective $l$ and hence enhances the dissipation, leading to attenuated amplitudes. The rightmost panel shows the data for KOI-54, which is qualitatively most consistent with shorter rotation periods of 1.0–2.0 d, comparable to the pseudo-synchronous period of $\sim 1.8$ d calculated in Section 5. (b) For edge-on systems, a similar argument regarding rotational suppression of mode amplitudes applies, but instead near the Doppler-shifted harmonic $k = 2P_{\text{orb}}/P_*$, rather than for $k$ near zero for the $m = 0$ modes observable face-on (see text for details). Comparison of the four left-hand panels to the rightmost panel, which shows our simple analytical equilibrium tide model’s harmonic decomposition (Appendix B2), demonstrates that the dynamical tide can dominate the light curve in edge-on systems.

Figure 8. Exploratory modelling of the Fourier power spectrum of KOI-54’s light curve. Our inhomogeneous theoretical model, including non-adiabaticity and rotation in the traditional approximation, is plotted above the graph’s horizontal axis, while the data for KOI-54 (Table 2) are plotted below. Parameters corresponding to this plot are in Table 4. Since harmonics 90 and 91 represent extreme resonances, they are difficult to resolve in a given stellar model grid, and fits which reproduce their amplitudes cannot reproduce other parts of the Fourier spectrum. Thus, we attempted to fit only $35 \leq k \leq 89$ and not the shaded region. Our fitting process was simplistic (see text), and we did not approach a full optimization, although our best fit does agree reasonably well with the data. Fig. 9 shows the observed flux perturbation from this plot for both stars separately on a fine frequency grid.

Also corresponds to an increase in the effective multipole $l$, where $l \sim \sqrt{\lambda}$, and $\lambda$ is a Hough eigenvalue from Section 3.4 (e.g. fig. 2 of Bildsten et al. 1996). Consequently, since $g$ modes asymptotically satisfy equation (6), the number of radial nodes $n$ must increase commensurately. Larger $n$ increases the radial wavenumber, which enhances the damping rate, further suppressing the resonant response of the modes and hence their contribution to the observed flux variation. This effectively corresponds to extending the highly damped travelling wave regime towards higher $k$ in Fig. 6.

As described in Section 3.4, when the magnitude of the Coriolis parameter becomes greater than unity, a new branch of eigenmodes develops with negative Hough eigenvalues, $\lambda < 0$. These modes are confined to the stellar poles rather than the equator (Lindzen 1966). They also have an imaginary Lamb frequency, so that they are radially evanescent (explained further in Fig. 1), and couple weakly to the tidal potential. We found negative-$\lambda$ modes to produce only a small contribution to the stellar response, which increased with increasing rotation rate but which was roughly constant as a function of forcing frequency, thus mimicking the equilibrium tide. The role of these modes in the context of tidal asteroseismology should be investigated further, but for now we have neglected them in Fig. 7.

For edge-on orbits, as in Fig. 7(b), the situation is more complicated, and there are high-amplitude pulsations observable at all rotation periods. First, $m = 0$ modes very weakly affect edge-on light curves, since their Hansen coefficients (which peak at $k = 0$) do not intersect with the linear overlap integrals as strongly as for
m = 2 modes (explained further in Section 4.2 and shown in Fig. 2). Similarly, modes with m = −2 have Hansen coefficients which peak near −2Ωorb/Ω1, and are very small for k ≥ 0.6. Thus, regardless of rotation, only m = +2 modes make significant light-curve contributions.

Within the m = +2 modes, there are two regimes to consider: prograde modes excited by harmonics k > 2Ω1/Ωorb and retrograde modes with k < 2Ω1/Ωorb (see Appendix C1). Prograde modes at a given corotating frame frequency are Doppler-shifted towards large k, whereas the Hansen coefficients peak near 2Ωperi/Ωorb, so their contribution to light curves is marginalized for fast rotation.

Retrograde, m = +2 g modes with small corotating-frame frequencies σlow = kΩorb − mΩ2 are subject to the same effect described earlier in the face-on case: they are suppressed by fast rotation due to weaker angular tidal coupling and stronger damping. The difference, however, is that although small driving frequencies are equivalent to small values of k for m = 0 modes, the Doppler shift experienced by m = 2 modes means that rotational suppression instead occurs for k ∼ 2Ω1/Ωorb, which is 84 × (P1/d) for KOI-54’s orbital period of 42 d. Fig. 7(b) demonstrates this, where e.g. little power can be observed near k = 84 for P1 = 1 d.

Furthermore, rotational suppression does not act on low-k harmonics in edge-on systems, as Fig. 7(b) also shows. Indeed, since fast rotation Doppler shifts lower retrograde modes − which radially couple more strongly to the tidal potential − towards values of k nearer to the Hansen peak of ∼2Ωperi/Ωorb, the power in lower harmonics can even be enhanced by sufficiently fast rotation rates.

The rightmost panel of Fig. 7(b) shows the harmonic decomposition of our simple equilibrium tide model for an edge-on orientation, not including irradiation (Section 6.1; Appendix B2). Comparing this plot to the left four panels shows in particular that, in edge-on orbits, the dynamical tide is not rotationally suppressed for harmonics where the equilibrium tide has large amplitudes, unlike for face-on orientations. Thus, the ellipsoidal variation of edge-on systems may be buried beneath the dynamical tidal response. This implies that full dynamical modelling may be necessary to constrain system parameters for edge-on binaries, and that care must be taken in searches for eccentric binaries, since it cannot be assumed that their light curves will be dominated by ellipsoidal modulations.

64 Light-curve power spectrum modelling

We performed preliminary quantitative modelling of the pulsation data in Table 2. As noted before, tidally driven pulsations should have frequencies which are pure harmonics of the orbital frequency, ω = kΩorb for k ∈ Z. Although most of the pulsations W11 report are of this form, some clearly are not, and are as such unaccounted for in linear perturbation theory. Hence we only attempted to model pulsations within 0.03 in k of a harmonic (set arbitrarily); this limited our sample to 21 harmonics, as shown in Table 2. We provide an explanation for the non-harmonic pulsations in Section 6.5.

There are eight primary parameters entering into our modelling of the remaining observed harmonics: stellar masses M1,2, radii R1,2, ZAMS metallicities Z1,2 and rigid-body rotation periods P1,2. To explore a range of stellar parameters, we used the stellar evolution code MESA (Paxton et al. 2011) to create two large sets of stellar models, one for each star, with ranges in M and R determined by W11’s constraints (Table 1). We set both stars’ metallicities to 0.04. The other two parameters, P1 and P2, were treated within our non-adiabatic code using the traditional approximation. We set P1 = P2 = 1.5 d, comparable to the expected pseudo-synchronous rotation period (Section 5.1) and qualitatively consistent with the small-amplitude flux perturbations of lower harmonics seen in KOI-54 (Section 6.3). We fixed all of the orbital parameters to the mean values given in W11.

As discussed in Section 6.6, it is possible that the 90th and 91st harmonics observed in KOI-54 are m = ±2 modes responsible for resonance locks, and are thus in states of nearly perfect resonance. Indeed, even if they are m = 0 chance resonances, which are ∼200 times easier to observe with KOI-54’s face-on orbital inclination than m = ±2 modes (Section 3.3), we find that a detuning of |δω/Ωorb| ∼ 10−2 is required to reproduce the amplitude of either harmonic, where δω is the difference between the eigenmode and driving frequencies (with δω = 0 representing a perfect resonance).

Such a close resonance represents a precise eigenfrequency measurement, and should place stringent constraints on stellar parameters. However, this degree of resonance is also very difficult to capture in a grid of stellar models because even changes in (say) mass of ΔM/Δρ ∼ 10−4 can alter the mode frequencies enough to significantly change the degree of resonance; future alternative modelling approaches may obviate this difficulty (Section 7). A second problem with trying to directly model the 90th and 91st harmonics is that the amplitudes of both of these harmonics may be set by non-linear processes, as addressed in Section 6.5. If correct, this implies that these particular modes strictly cannot be modelled using the linear methods we focus on in this paper.

We are thus justified in restricting our analysis to only those integral harmonics in the range 35 ≤ k ≤ 89. We chose k = 35 as our lower bound to avoid modelling harmonics that contribute to ellipsoidal variation. We set m = 0 for all of our analysis for the reason stated above. We also only used l = 2 for the tidal potential, since additional l terms are suppressed by (i) further powers of RP/Ω1, (ii) smaller disc-integral factors from Section 3.3 (e.g. b/2b2 = 0.2).

To find a reasonable fit to the harmonic power observed in KOI-54, we attempted a simplistic, brute-force optimization of our model against the data: we first modelled the linear response of each stellar model in our grid separately, ignoring its companion, and calculated the resulting observed flux perturbations as a function of k. We then compared the absolute values of these flux perturbations to the observations of KOI-54 and selected the best 10 parameter choices (M, R) for each star. (In future work, pulsation phases should be modelled in addition to the amplitudes reported by W11, since this doubles the information content of the data; see also further discussion of phases in Section 4.2.) Given this restricted set of stellar models, we computed the theoretical Fourier spectra for all 10 possible pairings of models.

Fig. 8 shows one of our best fits to the observations of KOI-54; Table 4 gives the associated stellar parameters. We obtained many reasonable fits similar to Fig. 8 with dissimilar stellar parameters, demonstrating that many local minima exist in this optimization.
problem. As a result, Fig. 8 and Table 4 should not be interpreted as true best fits but rather as an example of a model that can semi-quantitatively explain the observed harmonic power in KOI-54. We leave the task of using the observed pulsation data to quantitatively constrain the structure of the stars in KOI-54 to future work, as we discuss in Section 7.

Responses from both stars were used to create the plot in Fig. 8. Fig. 9, on the other hand, uses the same parameters, but instead shows each star’s observed flux perturbation separately and evaluated on a fine grid in frequency rather than only at integral orbital harmonics. As a result, Fig. 9 exposes the position of normal modes (which correspond to peaks in the black curves) in relation to observed harmonics (shown as red vertical lines), as well as other features not captured in the raw spectrum of Fig. 8. Fig. 9 also shows that harmonics 90 and 91 must come from different stars if they are indeed $m = 0$ g modes (although this may not be the case; see Section 6.6), since the g-mode spacing near $k \sim 90$ is much larger than the orbital period (with the same logic applying for harmonics 71 and 72).

6.5 Non-harmonic pulsations: three-mode coupling

W11 report nine pulsations which are not obvious harmonics of the orbital frequency; these have asterisks next to them in Table 2. As we showed previously (Section 4.2 and Appendix A1), these cannot be linearly driven modes. Here we present one possible explanation for the excitation of these pulsations.

To begin, we point out the following curious fact: the two highest-amplitude non-harmonic pulsations in Table 2 (F5 and F6) have frequencies which sum to 91.00 in units of the orbital frequency – precisely the harmonic with the second-largest amplitude (F2). (This is the only such instance, as we discuss below.)

Although this occurrence could be a numerical coincidence, it is strongly suggestive of parametric decay by non-linear three-mode coupling, the essential features of which we now describe. First, however, we emphasize that the treatment we present here is only approximate. In reality, the process of non-linear saturation is much more complicated, and a more complete calculation would involve fully coupling a large number of eigenmodes simultaneously (Weinberg & Quataert 2008).

If a parent eigenmode is linearly excited by the tidal potential to an amplitude that surpasses its three-mode-coupling threshold amplitude $S_\kappa$, any energy fed into it above that value will be bled away into daughter mode pairs, each with frequencies that sum to the parent’s oscillation frequency (Weinberg et al. 2011). In a tidally driven system, the sum of the daughter modes’ frequencies must thus be a harmonic of the orbital frequency.

For a parent with indices $a = (n, l, m)$ linearly driven at a frequency $\sigma$, the threshold is given by

$$|S_\kappa|^2 \simeq \min_{\kappa} \left( T_{abc} \right),$$

where

$$T_{abc} = \frac{\gamma_b \gamma_c}{4 \omega_a \omega_b \kappa_{abc}} \left( 1 + \frac{\delta \omega_{abc}^2}{(\gamma_a + \gamma_b)} \right).$$

$\omega_i$ is a mode frequency, $\gamma_i$ is a mode damping rate, $\kappa_{abc}$ is the normalization-dependent non-linear coupling coefficient (Schenk et al. 2002), $\delta \omega_{abc} = \sigma - \omega_b - \omega_c$ is the detuning frequency, and the minimization is over all possible daughter eigenmodes $b$ and $c$ (each short for an $(n, l, m)$ triplet). The non-linear coupling coefficient $\kappa_{abc}$ is non-zero only when the selection rules

$$0 = \text{mod}(l_a + l_b + l_c, 2),$$

$$0 = m_a + m_b + m_c,$$

$$|l_b - l_c| < l_a < l_b + l_c,$$

are satisfied. Due to the second of these rules, any Doppler shifts due to rotation do not affect the detuning since they must cancel.

For a simple system of three modes, the non-linear coupling’s saturation can be determined analytically. The parent saturates at

$$|S_\kappa|^2 \simeq \min_{\kappa} \left( T_{abc} \right),$$

where

$$T_{abc} = \frac{\gamma_b \gamma_c}{4 \omega_a \omega_b \kappa_{abc}} \left( 1 + \frac{\delta \omega_{abc}^2}{(\gamma_a + \gamma_b)} \right).$$

$\omega_i$ is a mode frequency, $\gamma_i$ is a mode damping rate, $\kappa_{abc}$ is the normalization-dependent non-linear coupling coefficient (Schenk et al. 2002), $\delta \omega_{abc} = \sigma - \omega_b - \omega_c$ is the detuning frequency, and the minimization is over all possible daughter eigenmodes $b$ and $c$ (each short for an $(n, l, m)$ triplet). The non-linear coupling coefficient $\kappa_{abc}$ is non-zero only when the selection rules

$$0 = \text{mod}(l_a + l_b + l_c, 2),$$

$$0 = m_a + m_b + m_c,$$

$$|l_b - l_c| < l_a < l_b + l_c,$$

are satisfied. Due to the second of these rules, any Doppler shifts due to rotation do not affect the detuning since they must cancel.

For a simple system of three modes, the non-linear coupling’s saturation can be determined analytically. The parent saturates at

$^7$This section uses the normalization of Weinberg et al. (2011), whereas the rest of the paper uses the normalization given in Section 3.2. We of course account for this when giving observable quantities.
the threshold amplitude $S_c$, and the ratio of daughter energies within each pair is given by the ratio of the daughters’ quality factors:

$$\frac{E_b}{E_c} = \frac{q_b}{q_c} = \frac{\omega_b/\gamma_b}{\omega_c/\gamma_c}. \quad (42)$$

Equations (37) and (38) exhibit a competition that determines which daughter pair will allow for the lowest threshold. At larger daughter $l$, modes are more finely spaced in frequency, since g-mode frequencies roughly satisfy the asymptotic scaling from equation (6); hence, the detuning $\delta \omega_{bc}$ becomes smaller (statistically) with increasing $l$. However, higher daughter $l$ also leads to increased damping rates at fixed frequency (equation 17). As such, the minimum threshold will occur at a balance between these two effects.

In order to semi-quantitatively address the phenomenon of three-mode coupling in KOI-54, we produced an example calculation of $S_c$, together with a list of best-coupled daughter pairs. To this end, we used a MESA stellar model (Paxton et al. 2011) consistent with the mean values of star 1’s properties reported in W11 (Table 1). We computed this model’s adiabatic normal modes using the ADIPLS code (Christensen-Dalsgaard 2008), and calculated each mode’s global quasi-adiabatic damping rate $\gamma_{\nu l}$ due to radiative diffusion (Section 3.2).

We focus on the second-highest-amplitude $k = 91$ harmonic present in the data (F1 from Table 2) and set $\sigma = 91 \times \Omega_{bc}$; as pointed out in W11, for $m_c = 0$ the quadrupolar eigenmode with natural frequency closest to the 91st orbital harmonic is the $g_{14}$ mode, i.e. the g mode with 14 radial nodes. We thus take this as our parent mode.

The minimization in equation (37) is over all normal modes, of which there is an infinite number. To make this problem tractable numerically, we essentially followed the procedure described in Weinberg & Quataert (2008):

(i) We restricted daughter modes to $1 \leq l \leq 6$. There is no reason a priori to suggest $l$ should be in this range, but, as shown in Table 5, $1 \leq l \leq 3$ turns out to be the optimum range for minimization in this particular situation, and modes with $l > 6$ are irrelevant.

(ii) The quantity to be minimized in equation (37), $T_{abc}$, achieves its minimum at fixed $\delta \omega_{bc}$ and $\kappa_{abc}$, for $\omega_b \approx \omega_c$, given the scaling from equation (17). As such, we computed all normal modes $b$ with frequencies in the range $f < \omega_b/\omega_a < 1 - f$; we took $f = 1/10$, which yielded 344 479 potential pairs, but trying $f = 1/5$, which yielded 61 623, did not change the result.

(iii) We computed $T_{abc}$, not including the three-mode-coupling coefficient $\kappa_{abc}$ (since it is computationally expensive to evaluate), for all possible pairs of modes satisfying (i) and (ii) as well as the selection rules in equations (39)–(41).

(iv) From the results of (iii), we selected the $N = 5000$ smallest threshold energies, and then recomputed $T_{abc}$ for these pairs this time including $\kappa_{abc}$ (Weinberg et al. 2011). (Trying $N = 1000$ did not change the results.) We set $m_b = m_c = 0 = m_e$ for simplicity, since $\kappa_{abc}$ depends only weakly on the values of $m$ so long as equation (40) is satisfied. Sorting again then yielded the best-coupled daughter pairs and an approximation for the saturation amplitude $S_c$.

Table 5 shows the best-coupled daughter mode pairs resulting from this procedure. It is interesting to note that most daughter pairs (i) involve an $l = 1$ mode coupled to an $l = 3$ mode (P2, P4, P5, P6) and/or (ii) have a large quality-factor ratio (all except P2, P5 and P9 have $|\frac{1}{2} \log_{10}(q_b/q_c)| > 0.5$).

For daughter pairs satisfying (i), the $l = 3$ mode would be much harder to observe in a light curve since disc averaging involves strong cancellation for larger-$l$ modes — indeed, Table 3 shows $b_1/b_9 \approx 10$ for Eddington limb darkening, where $b_1$ is a disc-integral factor defined in equation (23). The other disc-integral factor, $c_1$, does not decline as sharply with increasing $l$, but corresponds to cross-section perturbations, which are small relative to emitted flux perturbations as discussed above.) For daughter pairs satisfying (ii), since the ratio of daughter amplitudes scales as the square root of the ratio of their quality factors, one of the modes would again be difficult to observe.

Furthermore, if the parent had $m_b = \pm 2$ instead of $m_b = 0$ (see Section 6.6), each daughter pair would have several options for $m_b$ and $m_c$, introducing the possibility of $|m_b| \neq |m_c|$. This would mean daughters would experience even greater disparity in disc-integral cancellation due to the presence of $Y_{bc}^4(\theta, \phi_c)$ in equation (22); e.g. $|Y_{bc}^4(\theta, \phi_c) Y_{bc}^2(\theta, \phi_b)| \approx 0.02$ for KOI-54.

The above results provide a reasonable explanation for why there is only one instance of two non-harmonic pulsations adding up to an observed harmonic in the data for KOI-54 – only P9 from Table 5 has the potential to mimic pulsations F5 and F6 from W11. Nonetheless, the non-linear interpretation of the non-harmonic pulsations in KOI-54 predicts that every non-harmonic pulsation should be paired with a lower amplitude sister such that their two frequencies sum to an exact harmonic of the orbital frequency. This prediction may be testable given a sufficient signal-to-noise ratio, which may be possible with further observations of KOI-54.

Table 5. 10 best-coupled daughter mode pairs resulting from the procedure outlined in steps (i–iv) of Section 6.5. This is an example calculation and is not meant to quantitatively predict the non-harmonic components of KOI-54’s light curve. All frequencies and damping rates are in units of $\Omega_{bc}$. The square root of the daughter quality factor ratio, $\sqrt{q_b/q_c}$, gives an estimate of the ratio of daughter mode amplitudes, and hence of their potential relative light-curve contributions.

| ID | $(l_b, n_b)$ : $(l_c, n_c)$ | $\omega_b$ | $\alpha_b$ | $\alpha_c$ | $\frac{1}{2} \log_{10}(q_b/q_c)$ | $\log_{10}(|\delta \omega_{bc}|)$ | $\log_{10}(2\sqrt{q_b/q_c})$ |
|---|---|---|---|---|---|---|---|
| P1 | (2, −37) : (2, −23) | 35.3 | 55.8 | 0.66 | −1.4 | −2.1 |
| P2 | (1, −25) : (3, −30) | 29.9 | 61.0 | 0.057 | −1.5 | −2.5 |
| P3 | (1, −28) : (1, −11) | 26.8 | 64.1 | 0.69 | −1.3 | −3.1 |
| P4 | (1, −50) : (3, −24) | 15.3 | 75.7 | 1.2 | −1.4 | −1.7 |
| P5 | (1, −27) : (3, −29) | 27.8 | 63.2 | 0.031 | −1.4 | −2.5 |
| P6 | (1, −42) : (3, −25) | 18.0 | 73.1 | 1.2 | −1.6 | −1.7 |
| P7 | (1, −36) : (1, −10) | 20.9 | 69.8 | 1.3 | −0.61 | −2.7 |
| P8 | (2, −35) : (2, −24) | 37.2 | 53.7 | 0.51 | −0.87 | −2.2 |
| P9 | (2, −29) : (2, −28) | 44.7 | 46.4 | 0.038 | −0.99 | −2.5 |
| P10 | (1, −35) : (1, −10) | 21.5 | 69.8 | 1.2 | −0.51 | −2.8 |
Lastly, we can attempt to translate our estimate of the parent threshold amplitude $S_0$ into an observed flux perturbation, $\delta J_{cs}/J$, using the techniques of Section 3.3. Since our non-linear saturation calculation was performed with adiabatic normal modes, we strictly can only calculate the observed flux variation due to cross-section perturbations, $\delta J_{cs}$ (the $\xi_i$ component of equation 22), and not that due to emitted flux perturbations, $\delta J_{ef}$ (the $\Delta F$ component of equation 22). It evaluates to

$$\frac{\delta J_{cs}}{J} = |S_0 \times (2b_l - c_i) \times \xi_{cs}(R) \times Y_{20}(\theta_0, \phi_0)| \simeq 1.7 \text{ mmag}.$$ 

However, we can employ our non-adiabatic code to calibrate the ratio of $\delta J_{ef}$ to $\delta J_{cs}$, which we find to be $\delta J_{ef}/\delta J_{cs} \simeq 9$ for the 91st harmonic. We can then estimate the total saturated flux perturbation:

$$\frac{\delta J_{sat}}{J} = \left( \frac{\delta J_{ef}}{\delta J_{cs}} + 1 \right) \frac{\delta J_{cs}}{J} \simeq 17 \text{ mmag}.$$ 

This result is a factor of $\sim 100$ too large relative to the observed amplitude of 229 $\mu$mag for the 91st harmonic (Table 2). Taken at face value, this would mean that the inferred mode amplitude is below threshold, and should not be subject to non-linear processes, despite evidence to the contrary. There are several possible explanations for this discrepancy. If the 91st harmonic is actually an $m = \pm 2$ mode, which we proposed in Section 5.1, then the intrinsic amplitude required to produce a given observed flux perturbation is a factor of $\sim 200$ times larger than for $m = 0$ modes given KOI-54’s face-on inclination (Section 3.3). This would make the observed flux perturbation of the 91st harmonic comparable to that corresponding to the threshold for three-mode coupling, consistent with the existence of non-harmonic pulsations in the light curve. We discuss this further in Section 6.6.

Alternatively, if the 91st harmonic is in fact an $m = 0$ mode, many daughter modes may coherently contribute to the parametric resonance, reducing the threshold considerably, as in Weinberg et al. (2011). A more detailed calculation, coupling many relevant daughter and potentially granddaughte pairs simultaneously, should be able to address this more quantitatively.

### 6.6 Are harmonics 90 and 91 caused by prograde, resonance-locking, $|m| = 2$ g modes?

As introduced in Section 5.1, having two pseudo-synchronized stars presents an ostensibly appealing explanation for the large-amplitude 90th and 91st harmonics observed in KOI-54 (henceforth F1 and F2; Table 2): each is the manifestation of a different highly resonant eigenmode effecting a resonance lock for its respective star by opposing the equilibrium tide’s torque.

We discuss the viability of this interpretation below. First, however, what alternate explanation is available? The most plausible would be that F1 and F2 are completely independent, resonantly excited $m = 0$ modes. Each coincidence would require a detuning of $|\omega_{0d} - \sigma_{kw}|/\Omega_{orb} \sim 2 \times 10^{-2}$ (Fig. 9), which is equivalent to $|\omega_{0d} - \sigma_{kw}|/\Omega_{orb} \sim 10^{-4}$, where $\omega_{0d}$ is the nearest eigenfrequency and $\sigma_{kw} = \Omega_{orb} - m\Omega_s$ is the driving frequency. The probability of having a detuning equal to or smaller than this value, given $\sim 10$ available modes (Fig. 9), is $\sim 10$ per cent, so the combined probability if the resonances are independent is $\sim 1$ per cent. Moreover, in Section 6.4 we show that in this $m = 0$ interpretation, F1 and F2 must come from different stars, yet there is no explanation for why the two excited modes are so similar.

If instead F1 and F2 are due to highly resonant $m = \pm 2$ resonance locking modes, several observations are naturally explained. The high degree of resonance is an essential feature of the inevitable pseudo-synchronous state reached when the torque due to the dynamical tide cancels that due to the equilibrium tide (Section 5.1). The fact that the resonant modes correspond to similar $k$ would be largely a consequence of the fact that the two stars in the KOI-54 system are similar in mass and radius to $\sim 10$ per cent, so that a similar mode produces the dynamical tide torque in each star (although a corresponding $\sim 10$ per cent difference in $k$ would be equally possible in this interpretation).

In addition, we showed in Section 6.5 that the observed amplitudes of F1 and F2 are a factor of up to $\sim 100$ smaller than their non-linear threshold values assuming $m = 0$. There is also strong evidence that at least F2 has its amplitude set by non-linear saturation. Having $m \neq 0$ would help to resolve this discrepancy because the intrinsic amplitude of $m = \pm 2$ modes would need to be $\sim 200$ times larger to produce the observed flux perturbation. This would then imply that the amplitudes of F1 and F2 are indeed above the threshold for three-mode coupling, naturally explaining the presence of the non-harmonic pulsations in the KOI-54 light curve.

However, several significant problems with the resonance-locking interpretation arise upon closer examination. Assume that F1 and F2 indeed correspond to $m = \pm 2$ g modes that generate large torques effecting $P_{ps} \sim 1.8$ d pseudo-synchronization locks. In order to create positive torques, equation (C4) shows that we must have $m(k\Omega_{orb} - m\Omega_s) > 0$, which reduces to

$$k/m|\Omega_{orb} > \Omega_s.$$ 

In order to determine which modes correspond to F1 and F2, we can enforce a close resonance by setting

$$\omega_{0d} \simeq 90 \Omega_{orb} - 2 \Omega_s,$$

where we have used the fact that equation (43) requires $k$ and $m$ to have the same sign for a positive torque. For $l = 2$ and using a MESA stellar model consistent with W11’s mean modelled parameters for star 1 (Table 1), equation (44) yields $n \simeq 30$, neglecting rotational modification of the modes (i.e. not employing the traditional approximation).

However, in our calculations in Section 5.1 we find that the resonant torque due to the dynamical tide is instead typically caused by g modes with $n = 8-15$ (basically set by the intersection of the Hansen coefficient and linear overlap curves, as discussed in Section 4.2 in the context of flux perturbations). Using equation (44) again, this would mean we would expect $k \sim 140-200$. Furthermore, we find that even a perfectly resonant $n = 30$ g mode makes a negligible contribution to the torque. This is true both for ZAMS models and for evolved models consistent with the observed radii in KOI-54, indicating that there is little uncertainty introduced by the details of the stellar model. This result suggests that the g modes inferred to correspond to F1 and F2 are inconsistent with what would be expected from our torque calculation if the rotation rate is indeed $\sim 1.8$ d.

If we account for rotation in the traditional approximation (Section 3.4), the $n$ of a prograde mode of a given frequency can be at most a factor of $\sqrt{4/7}$ times smaller than its corresponding non-rotating value; this follows from the fact that the angular eigenvalue $\lambda$ asymptotes to $m^2 = 4$ in the limit $\omega_{0d} \ll \Omega_s$ for prograde modes, instead of $\lambda = (l + 1)/6$ in the non-rotating limit. This reduces the $n$ of the 90th harmonic from $n \simeq 30$ to $24$, still insufficient to yield a significant torque.
Another major problem with the resonance lock interpretation is that although our orbital evolution simulations described in Section 5.1 ubiquitously produce resonance locks, they always occur in only one star at a time. This is because if a mode is in a resonance lock in one star and a mode in the other star tries to simultaneously resonance lock, the first lock typically breaks since the orbital frequency begins to evolve too quickly for the lock to persist. Although it is possible for simultaneous resonance locks in both stars to occur, such a state is very improbable. Similar orbital evolution simulations presented in Fuller & Lai (2011) did produce simultaneous resonance locks, but only because they simulated only one star and simply doubled the energy deposition rate and torque, thus not allowing for the effect just described.

Finally, we point out one last inconsistency in the resonance lock interpretation of F1 and F2. It is straightforward to calculate the predicted flux perturbation associated with perfectly resonant \( |m| = 2 \) g modes in resonance locks (using e.g. the calibration discussed at the end of Section 6.5): for modes ranging from \( n \sim 8 \) to 15, we find that the predicted flux perturbation for KOI-54’s parameters is \( \sim 10\text{--}30 \mu\text{mag} \). This is a factor of \( \sim 10 \) smaller than the observed flux perturbations, yet somewhat larger than the smallest-amplitude pulsation reported by W11. It is also a factor of \( \sim 2 \) smaller than the non-linear coupling threshold for an \( m = 2 \) mode (which we determined using the same procedure as in Section 6.5, extended to allow for an \( m \neq 0 \) parent), although the uncertainties involved in our non-linear estimates are significant enough that we do not consider this to be a substantial problem.

Thus, even if F1 and F2 can be attributed to modes undergoing resonance locks (which is highly non-trivial, as we have seen), the observed amplitudes are larger than those we predict. Conversely, if F1 and F2 are simply chance \( m = 0 \) resonances, it appears that if a resonance lock existed, it would have been detected, although the possibility exists that the resonant mode’s flux perturbation was marginally smaller than those of the 30 reported pulsations due to uncertainties in our calculations. Firmer constraints on the flux perturbations in KOI-54 at \( k \sim 140\text{--}200 \) would be very valuable in constraining the existence of such \( m = \pm 2 \) modes, as would information about the phases of the 90th and 91st harmonics (see Section 4.2).

7 DISCUSSION

We have developed a set of theoretical tools for understanding and modelling photometric observations of eccentric stellar binaries. This work is motivated by the phenomenal photometry of the Kepler satellite and, in particular, by the discovery of the remarkable eccentric binary system KOI-54 (W11). This system consists of two similar A stars exhibiting strong ellipsoidal light-curve variation near periastron passage due to the system’s large \( (e = 0.83) \) eccentricity. W11 successfully modelled this phenomenon, and also reported the detection of at least 30 distinct sinusoidal pulsations in KOI-54’s light curve (Section 2), \( \sim 20 \) at exact harmonics of the orbital frequency and another \( \sim 10 \) non-harmonic pulsations. Although our work has focused on modelling KOI-54, our methods and techniques are more general, and are applicable to other similar systems.

We developed a simple model of KOI-54’s periastron brightening, including both the irradiation and equilibrium tide components of this effect, which agrees at the \( \sim 20 \) per cent level with the results W11 obtained using a much more detailed simulation (Section 6.1). Our model may be useful for analysis of other eccentric stellar binaries, allowing determination of orbital and stellar parameters; its simplicity should enable it to be implemented in an automated search of Kepler data.

In Section 4 we used the adiabatic normal mode formalism (see Section 3.2 and e.g. Kumar et al. 1995; Christensen-Dalsgaard 2003) to establish a qualitative connection between the range of stellar modes excited in a given binary system and the system’s orbital properties. For more detailed quantitative modelling of the harmonic pulsation spectrum of a given binary system, we further developed the non-adiabatic, inhomogeneous tidal method from Pfahl et al. (2008) by including the Coriolis force in the traditional approximation (Section 6.2; Appendix A).

In Section 6.3 we used this method to show that fast rotation tends to suppress power in the lower harmonics of a face-on binary system’s light curve (Fig. 7). This can qualitatively explain why there is a scarcity of large-amplitude, lower harmonic pulsations in KOI-54’s light curve, relative to predictions for non-rotating stars (Fig. 3). We also showed in Section 6.3, however, that the dynamical tidal response may be much larger than ellipsoidal variation in edge-on binaries, unlike in KOI-54 (which has an inclination of \( i = 5.5 \); see Table 1). For such systems, simultaneous modelling of the dynamical and equilibrium tides may be required in order to constrain system properties.

Moreover, in Section 5 we showed that rapid rotation periods of \( \sim 1.8 \text{~d} \) are expected for the A stars in KOI-54, due to pseudo-synchronization with the orbital motion near periastron. This pseudo-synchronous rotation period is shorter than the value of 2.53 d assumed by W11. The latter value is appropriate if the only appreciable torque is that produced by the equilibrium tide (Appendix C2). Since resonantly excited stellar \( g \) modes can produce a torque comparable to that of the equilibrium tide, pseudo-synchronous rotation can occur at even shorter rotation periods (Fig. 4). This involves a stochastic equilibrium between prograde resonance locks and the equilibrium tide. These same rapid rotation periods (\( \sim 1.8 \text{~d} \)) yield predicted light-curve power spectra that are the most qualitatively consistent with the pulsation data for KOI-54 (Fig. 7).

In Section 6.4 we performed a preliminary optimization of our non-adiabatic model by comparing its results in detail to the Fourier decomposition of KOI-54’s light curve (Table 2). We searched over an extensive grid of stellar masses and radii, assuming a metallicity of twice solar and a rotation period of 1.5 d. We also set \( m = 0 \), since KOI-54’s nearly face-on orientation implies that this is the case for almost all of the pulsations we modelled (Section 3.3). The modelling challenge in tidal asteroseismology contrasts with that of standard asteroseismology in that (i) we must simultaneously model both stars, and (ii) pulsation amplitudes and phases contain the key information in our case, since we are considering a forced system, whereas pulsation frequencies constitute the data in traditional asteroseismology. Moreover, stellar rotation is sufficiently rapid in eccentric binaries that its effect on stellar \( g \) modes cannot be treated perturbatively.

Although our minimization procedure was quite simple, we were able to obtain stellar models with power spectra semi-quantitatively consistent with the observations of KOI-54 (Figs 8 and 9). The resulting model in Fig. 8 is not formally a good fit, but this is not surprising given that two of the key parameters (metallicity and rotation period) were not varied in our analysis. Moreover, in our preliminary optimization we found that there were many local minima that produced comparably good light curves.

As noted above, a priori calculations suggest that both stars in the KOI-54 system should have achieved a pseudo-synchronous state at rotation periods of \( \sim 1.8 \text{~d} \). This requires frequent resonance locks to
occur, when a single $|m| = 2$ eigenmode comes into a near-perfect prograde resonance. A natural question is whether such a highly resonant mode could contribute to the KOI-54 light curve; this possibility is particularly attractive for the two largest-amplitude harmonics observed, the 90th and 91st. (See also our calculation of non-linear saturation from Section 6.5, discussed below.)

However, we find quantitative problems with this interpretation (Section 6.6). First, our orbital evolution simulations (Section 5.1) indicate that only one resonance lock should exist at a time, meaning that only one of the two large-amplitude harmonics could be explained in this way. This result is in disagreement with the simulations performed by Fuller & Lai (2011), since they did not simultaneously model both stars.

Further, in our calculations, the g modes capable of producing torques large enough to effect resonance locks have $n$ typically in the range 8–15 (where $n$ is the number of radial nodes), while the 90th harmonic corresponds to $n$ of 25–40 for $m = \pm 2$ and rotation periods of 2.0–1.5 d. Also, we predict that g modes producing resonance locks should have $k$ of 140–200, much larger than $\sim 90$, and flux perturbations of 10–30 µmag. The latter values are a factor of $\sim 10$ less than that observed for the 90th and 91st harmonics, but slightly larger than the smallest observed pulsations.

It thus seems quantitatively difficult to interpret harmonics 90 and 91 in KOI-54 as manifestations of $m = \pm 2$ modes in resonance locks, although we cannot conclusively rule out this possibility. Instead, it seems likely that they are simply chance $m = 0$ resonances (as is almost certainly the case for the overwhelming majority of the other observed pulsations in KOI-54). One theoretical uncertainty resides in our omission of rotational modification of the stellar eigenmodes when computing tidal torques. Our estimates suggest that this is a modest effect and is unlikely to qualitatively change our conclusions, but more detailed calculations are clearly warranted.

We note that in future work, pulsation phases should be modelled in addition to the amplitudes reported by W11, since this effectively doubles the information content of the data. Indeed, we showed in Section 4.2 that a resonant pulsation’s phase is strongly influenced by the mode’s value of $m$. In particular, since harmonics 90 and 91 are likely standing waves, as can be seen in the propagation diagram in Fig. 1, measurement of their phases could help to resolve the uncertainties pointed out above by supplying direct information about their degrees of resonance, thus potentially confirming or disproving the $m = \pm 2$ resonance lock interpretation.

In Section 6.5 we pointed out evidence for non-linear mode coupling in KOI-54’s observed pulsations: the existence of non-harmonic pulsations (which does not accord with linear theory; Section 4.2) and the fact that two of them have frequencies that sum to exactly the frequency of the 91st harmonic, the second-largest-amplitude harmonic pulsation in KOI-54’s light curve. This is consistent with parametric resonance, the leading-order non-linear correction to linear stellar oscillation theory (Weinberg et al. 2011).

Motivated by this observation, we performed a non-linear stability calculation that qualitatively explains why no other similar instance of a non-harmonic pair summing to an observed harmonic is present in the data: for the majority of daughter pairs likely to be non-linearly excited, there are sufficient differences in the $J$ and $m$ values of the daughter pair members, or sufficient differences in their predicted saturated energies, that only one member of the pair would be observable given current sensitivity. Nonetheless, the non-linear interpretation makes the strong prediction that every non-harmonic pulsation should be paired with a lower amplitude sister such that their two frequencies sum to an exact harmonic. This prediction may well be testable given a better signal-to-noise ratio.

One additional feature of the non-linear interpretation is that if the non-linearly unstable parent is an $m = 0$ mode, then the threshold amplitude for a linearly excited mode to be unstable to parametric resonance, which we have just argued exists in KOI-54, implies flux perturbations that are a factor of $\sim 100$ larger than those observed. In contrast, the parent being an $m = \pm 2$ mode ameliorates this discrepancy because the parent’s intrinsic amplitude must be $\sim 200$ times larger for a given flux perturbation due to KOI-54’s face-on orientation. This result thus argues in favour of the 91st harmonic in KOI-54 being an $m = \pm 2$ mode caught in a resonance lock, as discussed above.

There are many prospects for further development of the analysis begun in this paper. For example, in traditional asteroseismology, standard methods have been developed allowing a set of observed frequencies to be inverted uniquely, yielding direct constraints on stellar parameters, including the internal sound speed profile (Unno et al. 1989). The essential modelling difficulty in tidal asteroseismology is our inability to assign each observed pulsation amplitude to either star of a given binary a priori, hindering our attempts to develop a direct inversion technique. We leave the existence of such a technique as an open question.

Future observations of eccentric binaries may avoid this difficulty if one star is substantially more luminous than the other. However, for eccentric binaries with similar stars, in the absence of a means of direct inversion, we are left with a large parameter space over which to optimize, consisting a minimum of eight quantities: both stars’ masses, radii/ages, metallicities and rotation periods. Even this parameter set may ultimately prove insufficient, if modelling of tidally forced pulsations is found to be sensitive to the details of e.g. chemical mixing or convective overshoot, which can modify the Brunt–Väisälä frequency and thus g-mode frequencies.

One possible approach that should be explored in future work is to apply standard numerical optimization algorithms such as simulated annealing to this parameter space, attempting to minimize the $\chi^2$ of our non-adiabatic code’s theoretical Fourier spectrum against the observed harmonic pulsation data. In practice, it may be preferable to develop interpolation techniques over a grid of models given the high resolution in stellar parameters needed to resolve the close resonances responsible for large-amplitude pulsations.

Although KOI-54’s stars lie near the instability strip, this fact is unimportant for the tidal asteroseismology theory presented in this work. Consequently, future high-precision photometric observations of other eccentric binaries may supply a window into the structure of stars previously inaccessible by the techniques of asteroseismology. Constructing a data pipeline capable of reliably flagging eccentric binary candidates – e.g. finding efficient ways of searching for the equilibrium tidal/irradiation light-curve morphologies shown in Fig. 5 (Appendix B) – is also an important, complementary prospect for future work.

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The gravitational potential due to the secondary, experienced by the primary, is given by

\[ U_{2\rightarrow 1} = -\frac{GM_2}{|D - r|}. \]  

(A1)

Applying a multipole expansion (Jackson 1999) and excising the \( l = 0 \) (since it is constant) and \( l = 1 \) (since it is responsible for the Keplerian centre-of-mass motion) terms, we are left with the tidal potential:

\[ U = -\frac{GM_2}{D(t)} \sum_{l=2}^{\infty} \sum_{m=-l}^{l} W_{lm} \left( \frac{r}{D(t)} \right)^l e^{-i m f(t)} Y_{lm}(\theta, \phi), \]  

(A2)

where

\[ W_{lm} = \frac{4\pi}{2l + 1} Y_{lm}^* (\pi/2, 0) \]

\[ = (-1)^{l|m|} \text{mod}(l + m + 1, 2) \sqrt{\frac{4\pi}{2l + 1} \frac{(l + m - 1)!! (l - m - 1)!! (l + m)!! (l - m)!!}.} \]  

(A3)

Next, we shift to the primary’s corotating frame (by sending \( \phi \rightarrow \phi + \Omega_1 t \)) and expand the time dependence of the orbit in terms of the Hansen coefficients:

\[ U = \frac{M_2}{M_1} \sum_{l=0}^{\infty} \left( \frac{R_1}{a} \right)^{l+1} \sum_{m} W_{lm} Y_{lm}(\theta, \phi) \sum_{k} \exp(-i\sigma_{l,m} t) X_{lm}^k(r) U_l(r), \]  

(A4)

with \( \sigma_{l,m} = k\Omega_{\text{orb}} - m\Omega_\ast \) and

\[ U_l(r) = -\left( \frac{GM_1}{R_1} \right) \left( \frac{r}{R_1} \right)^l. \]  

(A5)

The unit-normalized Hansen coefficients \( \tilde{X}_{lm}^k \) were defined in equation (11); here we are using the conventionally normalized Hansen coefficients \( X_{lm}^k = \tilde{X}_{lm}^k/(1 - e)^{l+1} \), which are convenient to evaluate numerically as an integral over the eccentric anomaly:

\[ X_{lm}^k = \frac{1}{\pi} \int_0^\pi (1 - e \cos E)^{l+1} \cos \left( kE - e \sin E \right) - 2m \arctan \left( \sqrt{\frac{1 + e}{1 - e} \tan(E/2)} \right) dE. \]  

(A6)
If we represent the linear response of a star to the perturbing tidal potential by an abstract vector \( y(r, \theta, \phi, t) \) whose components are the various oscillation variables (e.g., \( \xi/r \)), then \( y \) can also be expanded, again in the primary’s corotating frame, as in (A4):

\[
y = \frac{M_2}{M_1} \sum \left( \frac{R_1}{a} \right)^{l+1} \sum_m W_{lm} Y_{lm}(\theta, \phi) \sum_k \exp(-i\sigma_{km} t) X^k_{lm}(\epsilon) y^k_{lm}(r). \tag{A7}
\]

The equations necessary to determine \( y^k_{lm}(r) \) are given in the appendix of Pfahl et al. (2008), along with appropriate boundary conditions; note that their \( U \) is our \( U_i \) and their driving frequency \( \omega \) is our \( \sigma_{km} \).

After determining \( y^k_{lm}(r) \) in the corotating frame, we can switch to the inertial frame specified in Section 3.1:

\[
y = \frac{M_2}{M_1} \sum \left( \frac{R_1}{a} \right)^{l+1} \sum_k \exp(-i\Omega_1 t) \sum_m W_{lm} Y_{lm}(\theta, \phi) X^k_{lm}(\epsilon) y^k_{lm}(r). \tag{A8}
\]

As noted in Section 4.2, we see in equation (A8) that the observed frequencies should be pure harmonics of the orbital frequency, even though the corresponding amplitudes of observed pulsations are influenced by the star’s rotation rate (via the Doppler-shifted frequency \( \sigma_{km} \)).

### A2 Rotation in the traditional approximation

We now invoke the traditional approximation (Section 3.4); we must correspondingly adopt the Cowling approximation and employ the Hough functions (Section 5) as angular basis functions instead of spherical harmonics.

We expand the Hough functions as (Longuet-Higgins 1968)

\[
H^k_{lm} = \sum_l e^k_{lm} \tilde{P}_l \rightarrow e^k_{lm} = 2\pi \int_0^1 \tilde{P}_l H^k_{lm} d\mu \rightarrow \tilde{P}_l = \sum_l e^k_{lm} H^k_{lm}, \tag{A9}
\]

where \( \tilde{P}_l \) is a normalized associated Legendre function defined by

\[
\tilde{P}_l = \frac{\sqrt{2l + 1} (l - m)!}{4\pi (l + m)!} P_{lm} \rightarrow Y_{lm}(\theta, \phi) = e^{i\theta} \tilde{P}_l (\cos \theta). \tag{A10}
\]

We used the numerical method of calculating the expansion coefficients \( e^k_{lm} \) detailed in Ogilvie & Lin (2004) Section 5.4. The tidal potential in the corotating frame is then

\[
U = \frac{M_2}{M_1} \sum \left( \frac{R_1}{a} \right)^{l+1} \sum_m W_{lm} e^{i\theta} \sum_k \exp(-i\sigma_{km} t) X^k_{lm} \sum_{\lambda} H^k_{\lambda m}(\mu) U^k_{\lambda m}(r), \tag{A11}
\]

where

\[
U^k_{\lambda m}(r) = - \left( \frac{GM_1}{R_1} \right) \left( \frac{r}{R_1} \right)^l e^k_{lm}, \tag{A12}
\]

\[\sigma_{km} = k\Omega_{orb} - m\Omega,\] and the Coriolis parameter \( q \) on which the Hough functions depend is

\[q = 2\Omega_1/\sigma_{km}; \tag{A13}\]

which justifies writing \( H^k_{\lambda m} \) and \( e^k_{lm} \), rather than \( H^q_{\lambda m} \) and \( e^q_{lm} \).

We again represent the linear response of a star, as in Appendix A1, by a vector \( y(r, \theta, \phi, t) \) whose components are the various oscillation variables, and which can be expressed in the inertial frame as

\[
y = \frac{M_2}{M_1} \sum_k \exp(-i\Omega_{orb} t) \sum_l \left( \frac{R_1}{a} \right)^{l+1} \sum_m W_{lm} e^{i\theta} X^k_{lm}(\epsilon) \sum_{\lambda} H^k_{\lambda m}(\mu) y^k_{\lambda m}(r) \tag{A14}
\]

The expansion of \( H^k_{\lambda m} \) back into associated Legendre functions in the second line of equation (A14) is useful since disc integrals are convenient to perform over spherical harmonics (Section 3.3).

Following Unno et al. (1989), we choose the components of \( y \) as

\[
y_1 = \frac{\xi_1}{r}, \quad y_2 = \frac{\delta p}{\rho gr}, \quad y_3 = \frac{\Delta s}{c_p}, \quad \text{and} \quad y_6 = \frac{\Delta L}{L_\ast}, \tag{A15}
\]

where we have omitted the variables corresponding to the perturbed gravitational potential, \( y_3 \) and \( y_4 \). Equation (A14) together with determination of the radial displacement \( \xi_1/r = y_1 \) and the Lagrangian flux perturbation \( \Delta F/F = y_6 - 2y_4 \) at the photosphere then enables the use of the formalism from Section 3.3 to compute the flux perturbation as seen by an observer.

Next, we present the differential equations which determine a particular component \( y^k_{\lambda m}(r) \) of the full response in radiative zones. These equations are nearly identical to those in the appendix of Pfahl et al. (2008), but with \( R_l^{l+1} \) replaced by \( l \) and with certain terms set to zero as per the traditional approximation. In practice these terms can be left in, since they are nearly zero for situations where the traditional
approximation is valid; this is then a smooth way of transitioning among different regimes. Omitting \( \lambda \cdot \ln k \) indices and denoting \( U = U_{\lambda m} \), and \( \omega = \sigma_{\lambda m} \), the equations are

\[
\frac{dy_1}{d \ln r} = y_1 \left( \frac{gr}{c_s^2} - 3 \right) + y_2 \left( \frac{\lambda g}{\omega^2 r} - \frac{gr}{c_s^2} \right) - y_2 \rho s + \frac{\lambda}{\omega^2 r^2} U, \\
\frac{dy_2}{d \ln r} = y_1 \left( \frac{\omega^2 - N_s^2}{g/r} \right) + y_2 \left( 1 - \eta + \frac{N_s^2}{g/r} \right) - y_3 \rho s - \frac{1}{g} \frac{dU}{dr}, \\
\frac{dy_3}{d \ln r} = y_3 \frac{r}{H_p} \left[ \nabla_{ad} \left( \eta - \frac{\omega^2}{g/r} \right) + 4(\nabla - \nabla_{ad}) + c_2 \right] + y_2 \frac{r}{H_p} \left[ \left( \nabla_{ad} - \nabla \right) \frac{\lambda g}{\omega^2 r} = c_2 \right] \\
+ y_3 \frac{r}{H_p} \nabla (4 - \kappa_s) - y_3 \frac{r}{H_p} \nabla + \frac{r}{H_p} \left[ \nabla_{ad} \left( \frac{dU/dr}{g} \right) + \left( \nabla_{ad} - \nabla \right) \frac{\lambda}{\omega^2 r^2} U \right], \\
\frac{dy_4}{d \ln r} = y_2 \left( \frac{\lambda g}{\omega^2 r} \right) + y_3 \left( \frac{4\pi r^3 \rho_c T}{L} \right) - y_4 \gamma + \left( \frac{\lambda \gamma}{\omega^2 r} \right) U,
\]

where \( \eta = 4\pi r^3 \rho /M_s, \gamma = 4\pi r \rho s /L_s, c_2 = (r/H_p) \nabla (\kappa_{ad} - 4\nabla_{ad}) + \nabla_{ad} (d \ln \nabla_{ad} /d \ln r + r/H_p), H_p = \rho g /p \) is the pressure scale height, and \( \epsilon \) is the specific energy generation rate.

We need four boundary conditions for our four variables. Our first three are

\[
0 = \xi_i(0) \quad \text{evanescence in convective core}, \\
0 = \Delta s(0) \quad \text{adiabaticity/evanescence in core}, \\
0 = \frac{\Delta F(R)}{F(R)} - 4 \frac{\Delta T(R)}{T(R)} \quad \text{blackbody at the stellar surface},
\]

where \( \Delta T / T \) can be cast in terms of \( y_1, y_2 \) and \( y_3 \) using standard thermodynamic derivative identities.

A final surface boundary condition that allows for travelling and/or standing waves can be generated by imposing energetic constraints at the surface. This is detailed in Unno et al. (1989) for adiabatic oscillations. To generalize the boundary condition to include non-adiabatic, rotation and inhomogeneous tidal forcing, we write equations (A16)–(A19) as

\[
\frac{dy}{d \ln r} = M y + b,
\]

where \( M \) and \( b \) are treated as constant near the stellar photosphere. The constant solution is \( y_0 = -M^{-1} b \); defining \( z = y - y_0 \), the homogeneous solutions for \( z \) can be computed by diagonalizing \( M \). In the evanescent case, we eliminate the solution for \( y \) with outwardly increasing energy density. Alternatively, in the travelling wave case, we eliminate the inadwar-propagating wave. The final boundary condition is then implemented by setting the amplitude of the eliminated homogeneous solution to zero, and solving for a relationship between the original fluid variables implied by this statement.

**APPENDIX B: ANALYTIC MODEL OF ELLIPSOIDAL VARIATION**

As discussed in Section 6.1, our simplified model of ellipsoidal variation reproduces the much more sophisticated simulation code employed by Welsh et al. (2011) to model KOI-54; here we discuss the details of our analytic methods, which can easily be applied to model other systems.

**B1 Irradiation**

The following is our simple analytical model of the insolation component of the KOI-54’s ellipsoidal variation. We focus our analysis on the primary, since extending our results to the secondary is trivial. Our main assumption is that all radiation from the secondary incident upon the primary is immediately reprocessed at the primary’s photosphere and emitted isotropically (i.e. absorption, thermalization and re-emission). This assumption is well justified for KOI-54, since its two component stars are of very similar spectral type. The method below might need to be modified if the components of a binary system had significantly different spectral types, because then some of the incident radiation might instead be scattered.

The incident flux on the primary, using the conventions and definitions introduced in Section 3.1, is

\[
F_{z \rightarrow 1} = \frac{L_z}{4\pi D^2} Z(\hat{r} \cdot \hat{D}),
\]

where \( Z \) is the ramp function, defined by

\[
Z(x) = \begin{cases} 
0 & x < 0 \\
1 & x \geq 0 
\end{cases}
\]

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We can expand \( Z(\hat{r} \cdot \hat{D}) \) in spherical harmonics as

\[
Z(\hat{r} \cdot \hat{D}) = \sum_{lm} Z_{lm} Y_{lm}(\theta, \phi) e^{-im\varphi},
\]

where \( Z_{lm} \) can be evaluated to

\[
Z_{lm} = 2 \left( \frac{2l + 1 (l - m)!}{4\pi (l + m)!} \right)^{1/2} \left( \frac{\cos(m\pi/2)}{1 - m^2} \right) \int_{-1}^{1} \sqrt{1 - \mu^2} P_{lm} d\mu,
\]

with \( \cos(m\pi/2)/(1 - m^2) \to \pi/2 \) for \( m = \pm 1 \).

Next, taking the re-emission as isotropic, the re-emitted intensity will be

\[
I_{\text{emit}} = \frac{F_2 - 1}{\pi}.
\]

Using this together with our expansion of \( Z(\hat{r} \cdot \hat{D}) \) as well as results from Section 3.3, we can evaluate the observed flux perturbation:

\[
\frac{\delta J}{J_1} = \beta(T_1) \left( \frac{L_2}{L_1} \frac{R_1^2}{D(t)} \right) \sum_{i=0}^{\infty} \sum_{m=-l}^{l} b_i Z_{lm} Y_{lm}(\theta_o, \phi_o) e^{-im\varphi},
\]

where \( J_1 = L_1/4\pi c^2 \) is the unperturbed observed flux, \( s \) is the distance to the observer, the bandpass correction coefficient \( \beta(T) \) is defined in equation (19), the disc-integral factor \( b_i \) is defined in equation (23), several values of \( b_i \) using Eddington limb darkening are given in Table 3, and other variables are defined in Section 3.1. Since \( b_i \) declines rapidly with increasing \( l \), it is acceptable to include only the first few terms of the sum in equation (B6). We have neglected limb darkening, so it is formally necessary to use a flat limb-darkening law in calculating \( b_i \) [\( b(\mu) = 2 \); Section 3.3]. However, we found this to be a very modest effect.

The binary separation \( D \) and the true anomaly \( f \) can be obtained as functions of time in various ways, e.g. by expanding with the Hansen coefficients employed earlier (equation 11 or A6), or by using

\[
D = a(1 - e^2) \frac{1}{1 + e \cos f},
\]

together with numerical inversion of

\[
\Omega_{\text{obs}} = 2 \arctan \left( \frac{1 - e \tan (f/2)}{\sqrt{1 - e^2}} \right) = \frac{e \sqrt{1 - e^2 \sin f}}{1 + e \cos f}, -\pi < f < \pi.
\]

The observed flux perturbation from the secondary is obtained from equation (B6) by switching \( 1 \leftrightarrow 2 \) and sending \( \phi_o \to \phi_o + \pi \).

Using the fact that \( b_0 = 1 \) and \( Z_{00} = \sqrt{T_2}/2 \), it can be readily verified that the total reflected power, i.e. \( s^2 J_1 \) times (B6) integrated over all observer angles (\( \theta_o, \phi_o \)), is equal to \( L_2(\pi R_1^2/4\pi D^2) \). This is just the secondary’s luminosity times the fraction of the secondary’s full solid angle occupied by the primary, which is the total amount of the secondary’s radiation incident on the primary.

### B2 Equilibrium tide

We invoke the Cowling approximation (well satisfied for surface values of perturbation variables), and use the analytic equilibrium tide solution, where the radial displacement at star 1’s surface becomes

\[
\xi = -\frac{U(R_1, t)}{g(R_1)},
\]

and \( U \) is the tidal potential. Using the expansion in equation (A2), \( \xi_{1, \text{eq}}(t)/R_1 \) from equation (20) becomes (Goldreich & Nicholson 1989)

\[
\xi(t) = \frac{M_2}{M_1} \left( R_1 \frac{D(t)}{D(t)} \right) \sum_{i=0}^{\infty} \sum_{m=-l}^{l} b_i Z_{lm} Y_{lm}(\theta_o, \phi_o) e^{-im\varphi},
\]

We invoke von Zeipel’s theorem (von Zeipel 1924; Pfahl et al. 2008) to determine the corresponding surface emitted flux perturbation:

\[
\frac{\Delta F_{\text{em}}(t)}{F_1} = -(l + 2) \frac{\xi(t)}{R_1}.
\]

We can then explicitly evaluate the observed flux variation using the formalism from Section 3.3:

\[
\frac{\delta J}{J_1} = \frac{M_2}{M_1} \sum_{i=0}^{\infty} \sum_{m=-l}^{l} \left[ (2 - \beta(T_1))(l + 2)b_i - c_i \right] W_{lm} Y_{lm}(\theta_o, \phi_o) e^{-im\varphi},
\]

where the bandpass correction coefficient \( \beta(T) \) is defined in equation (19), \( W_{lm} \) is defined in equation (A3), the disc-integral factors \( b_i \) and \( c_i \) are defined in equations (23) and (24), several values of \( b_i \) and \( c_i \) using Eddington limb darkening are given in Table 3, and other variables are defined in Section 3.1. Due to the strong dependence on \( l \), it is typically acceptable to include only the first term of the sum in equation (B6). Computation of the binary separation \( D(t) \) and true anomaly \( f(t) \) is discussed in Appendix B1. The observed flux perturbation from the secondary is obtained from equation (B12) by switching \( 1 \leftrightarrow 2 \) and sending \( \phi_o \to \phi_o + \pi \).

We note that although the analytic equilibrium tide solution for the radial displacement \( \xi \) is a good approximation at the stellar surface regardless of stellar parameters, the presence of a significant surface convection zone in a solar-type star proscribes the use of equation (B11); Pfahl et al. (2008) gives the appropriate replacement in their equation (37). Moreover, we note that equation (B11) may also be invalid for slowly rotating stars in eccentric orbits; see Section 6.2.

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APPENDIX C: TIDAL ORBITAL EVOLUTION

C1 Eigenmode expansion of tidal torque and energy deposition rate

Assuming alignment of rotational and orbital angular momenta, the tidal torque $\tau$ produced by star 2 on star 1 must have only a $z$-component, where $\hat{z}$ points along the orbital angular momentum. We can evaluate it as follows (Kumar & Quataert 1998). First,

$$\tau = \hat{z} \cdot \int_a^b \left( r \times \frac{dF}{dV} \right) dV = \int_a^b \left( \hat{z} \times r \right) \cdot \frac{dF}{dV} dV$$

$$= \int_a^b \hat{\phi} \cdot [-(\rho_0 + \delta\rho) \nabla U] r \sin \theta \, dV.$$  \hfill (C1)

The term involving the background density $\rho_0$ vanishes; expanding both the tidal potential $U_{lm}(r, t)$ and the Eulerian density perturbation $\delta\rho$ in spherical harmonics with expansion coefficients $\sigma_{lm}$ we have

$$\tau(t) = i \sum_{l=2}^{\infty} \sum_{m=-l}^{l} m \int_0^{R_i} \delta\rho_{lm}(r, t) U_{lm}^*(r, t) r^2 \, dr.$$  \hfill (C2)

Further invoking the expansions from equations (5) and (A4), as well as the definitions in Section 3.2, we arrive with

$$\tau(t) = -2i \left( \frac{GM_2^2}{R_1^2} \right) \sum_{nlm} m \left( \frac{Q_{nl} W_{lm}^2}{E_{nl}} \right) X_{lm}^k X_{lm}^* \Delta_{nlm} e^{i(k^2-4\Omega_{orb})t}.$$  \hfill (C3)

Lastly, averaging over a complete orbital period and rearranging the sums, we derive our final expression for the secular tidal torque:

$$\langle \tau \rangle = 8\Omega_{orb} \left( \frac{GM_2^2}{R_1} \right) \left( \frac{M_2}{M_1} \right)^2 \sum_{i=2}^{\infty} \left( \frac{R_i}{a} \right)^{2i+2} \sum_{m=-l}^{l} W_{lm}^2 \sum_{k=0}^{\infty} k X_{lm}^k (e^2)^2 \sum_n \left( \frac{Q_{nl}^2}{E_{nl}} \right) \left( \frac{\omega_{nl}^2 \gamma_{nl} Y_{nl}}{(\omega_{nl}^2 - \sigma_{lm}^2)^2 + 4\gamma_{nl}^2 \sigma_{lm}^2} \right).$$  \hfill (C4)

The torque depends on the rotation rate $\Omega$, only through the Doppler-shifted frequency $\sigma_{lm} = k \Omega_{orb} - m \Omega_\star$, since we have neglected rotational modification of the eigenmodes (Section 3.4). Fig. 4 shows plots of this torque evaluated for KOI-54.

Note that a particular term of this sum is positive if and only if $k \sigma_{lm} + m \omega_\star < 0$, which reduces to $(k/m)\Omega_{orb} > \Omega_\star$. This is known as being prograde, since it is equivalent to the condition that a mode’s angular structure, in the corotating frame, rotates in the same sense as the stellar spin; conversely, retrograde waves with $(k/m)\Omega_{orb} < \Omega_\star$ cause negative torques.

Using similar techniques to those given above, an equivalent expansion of the secular tidal energy deposition rate into the star (including mechanical rotational energy) can be derived:

$$\langle E \rangle = 8\Omega_{orb} \left( \frac{GM_2^2}{R_1} \right) \left( \frac{M_2}{M_1} \right)^2 \sum_{i=2}^{\infty} \left( \frac{R_i}{a} \right)^{2i+2} \sum_{m=-l}^{l} W_{lm}^2 \sum_{k=0}^{\infty} k X_{lm}^k (e^2)^2 \sum_n \left( \frac{Q_{nl}^2}{E_{nl}} \right) \left( \frac{\omega_{nl}^2 \gamma_{nl} Y_{nl}}{(\omega_{nl}^2 - \sigma_{lm}^2)^2 + 4\gamma_{nl}^2 \sigma_{lm}^2} \right).$$  \hfill (C5)

The only difference between equations (C4) and (C5) is switching $m \leftrightarrow k\Omega_{orb}$.

C2 Non-resonant pseudo-synchronization

A pseudo-synchronous frequency $\Omega_{ps}$ is defined as a rotation rate that produces no average tidal torque on the star throughout a sufficiently long time interval, which here we take to be a complete orbital period (Section 5), i.e.

$$\langle \tau \rangle (\Omega_{ps}) = 0.$$  \hfill (C6)

Here we will show that our expansion from (C1) reproduces the value of $\Omega_{ps}$ derived in Hut (1981), which we denote $\Omega_{ps}^\infty$, in the equilibrium tide limit. We will in particular show that Hut’s result is independent of assumptions about eigenmode damping rates.

Proceeding, we take the non-resonant (equilibrium tide) limit of equation (C4). This is obtained by retaining only the first term in the Taylor series expansion in $\sigma_{lm}/\omega_{nl}$ of the last factor in parentheses from equation (C4), and yields

$$\langle \tau_{nl} \rangle = 8 \left( \frac{GM_2^2}{R_1} \right) \left( \frac{M_2}{M_1} \right)^2 \sum_{i=2}^{\infty} \left( \frac{R_i}{a} \right)^{2i+2} \sum_{m=-l}^{l} W_{lm}^2 \sum_{k=0}^{\infty} k X_{lm}^k (e^2)^2 \sum_n \left( \frac{Q_{nl}^2}{E_{nl}} \right) \left( \frac{\omega_{nl}^2 \gamma_{nl} Y_{nl}}{(\omega_{nl}^2 - \sigma_{lm}^2)^2 + 4\gamma_{nl}^2 \sigma_{lm}^2} \right);$$  \hfill (C7)

note that sums over $k$ and $m$ become decoupled from the sum over $n$. Setting $\langle \tau_{nl} \rangle (\Omega_{ps}^\infty) = 0$ and retaining only $l = 2$, we have

$$0 = \sum_{k=\infty}^{\infty} X_{lm}^2 (e^2)^2 \left( k \Omega_{orb} - 2 \Omega_{ps}^\infty \right).$$  \hfill (C8)

We need two identities to evaluate this further. First, starting with the definition of the Hansen coefficients,

$$\left( \frac{a}{D} \right)^{l+1} e^{-imf} = \sum_{k=\infty}^{\infty} X_{lm}^k e^{-ik\Omega_{orb}t},$$  \hfill (C9)

we can differentiate with respect to $t$, then multiply by the complex conjugate of (C9) and average over a complete period to derive

$$\sum_{k=\infty}^{\infty} k (X_{lm}^k)^2 = \frac{m}{2\pi} \int_0^{2\pi} \left( \frac{1 + e \cos f}{1 - e^2} \right)^{2i+2} df.$$  \hfill (C10)

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Specializing to \( l = 2 \),
\[
\sum_{k=-\infty}^{\infty} k (X_{2m}^k)^2 = m \left[ \frac{5e^6 + 90e^4 + 120e^2 + 16}{16(1 - e^2)^6} \right].
\] (C11)

The second identity needed,
\[
\sum_{k=-\infty}^{\infty} (X_{2m}^k)^2 = \frac{3e^4 + 24e^2 + 8}{8(1 - e^2)^{3/2}},
\] (C12)

can be derived similarly.

Substituting equations (C11) and (C12) into (C8), we have that
\[
\Omega_{ps}^n = \Omega_{orb} \frac{1 + (15/2)e^2 + (45/8)e^4 + (5/16)e^6}{[1 + 3e^2 + (3/8)e^4](1 - e^2)^{3/2}},
\] (C13)

this is precisely equation (42) from Hut (1981).

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