

Examining Instabilities Due to Driven Scalars in AdS

ABSTRACT: We extend the study of holographic quantum quenches in AdS to include solutions to driven systems, i.e. those with time-dependent sources on the AdS boundary. This necessitates the activation of non-normalizable modes in the massive bulk scalar field, which then couple to the metric and normalizable scalar modes. Analytic expressions for secular terms that persist after time averaging are determined for scalars in AdS_{d+1} with any mass, and for different driving frequencies. We then numerically evaluate these sources for $d = 4$ and discuss what role these play in the perturbative stability of these systems.

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1 Introduction

Examinations of quenches in strongly coupled quantum systems can be approached using the holographic dual of such a process, namely the evolution of weakly coupled gravity. Full accounting of the return to equilibrium – signalled by the formation of a black hole in the dual theory – requires advanced numerical methods, and has been the subject of numerous studies. While AdS has been shown to be perturbatively unstable to generic data, an interesting array of stable and meta-stable behaviours that resist gravitational collapse have been demonstrated for initial data close to the fundamental AdS modes, the oscillons. Data that exhibits stability over long times (a scale that is often set by the amplitude of the perturbation) exhibits inverse energy cascades that balance the direct cascade of energy to short length scales. This weakly turbulent energy cascade is captured by the third-order dynamics of a perturbative expansion. To describe the balance of inverse and direct energy flow, a second, “slow time” is introduced that governs the evolution of the scalar field and, therefore, the metric functions. This is known as the Two-Time Formulation (TTF), and allows for analytical determination of the evolution of the scalar field in the perturbative regime.

Conventional examinations perturbative stability using TTF have focused on the quenches of some initial energy perturbation. However, the analytic descriptions of holographic pumped solutions – those with periodic boundary conditions that constantly inject energy into the system – remains undetermined. With this in mind, we examine the effect of a time-dependent source on the conformal boundary has on the analytic expressions for the time evolution of the slowly varying integration constants. Quenches in asymptotically AdS spacetime restrict the space of oscillon solutions to those that approach zero near the conformal boundary; however, in a driven solution the energy is pumped into the system through a second class of oscillons: those that approach constant, non-zero values on the boundary. Since these solutions will have non-finite inner products over the space, they are known as non-normalizable solutions. These non-normalizable modes couple to the metric and the normalizable modes to bring energy into the system, where direct and inverse energy cascades proceed over perturbative time scales.

Following the study of resummation and time-averaging procedures for scalar fields in AdS, we isolate secular terms – those that grow linearly with time and cannot be absorbed by a phase term – from those that are averaged out. The terms that persist after time averaging are those that obey certain resonance conditions between the frequencies of the non- and normalizable modes. By evaluating the third-order interactions on resonance, we use a renormalization procedure to absorb resonant contributions into the equations for the slowly varying amplitude and phase variables of the scalar field.

This paper is organized as follows: after a brief discussion of how to arrive at the third order source term, we consider the addition of a time-dependent boundary condition for the scalar field. As an exercise, and to provide explicit expressions for our choice of gauge and mass, §3 examines the resonant contributions in the case of a massive scalar field in AdS_{d+1} with any mass-squared, up to and including the Breitenlohner-Freedman mass [1]: $m_{BF}^2 \leq m^2$.

We recover the natural vanishing of two of the three resonances, and then examine the effects of mass-dependence on the non-vanishing channel. Whenever values are calculated, the choice of $d = 4$ is implied as to draw the most direct comparison to existing literature. In section §4, we extend the boundary conditions to include a variety of periodic boundary sources that couple to non-normalizable modes in the bulk. For each choice of boundary condition, we derive analytic expressions for applicable resonances and evaluate these expressions for different ranges of scalar field masses. Finally, in §5 we discuss the implications of non-vanishing resonances on the competing energy cascades and the perturbative stability of such systems. For completeness, we include details of our derivation of the general source term in an appendix, as well as a complete list of possible resonances and their contributions in the case of activating two, equal frequency non-normalizable modes.

2 Source Terms and Boundary Conditions

Let us first consider a minimally coupled, massive scalar field coupled to a spherically symmetric, asymptotically AdS_{d+1} spacetime in global coordinates, whose metric is given by

$$ds^2 = \frac{L^2}{\cos(x)} \left(-A(t, x) e^{-2\delta(t, x)} dt^2 + A^{-1}(t, x) dx^2 + \sin^2(x) d\Omega_{d-1}^2 \right), \quad (2.1)$$

where L is the AdS curvature (hereafter set to 1), and the radial coordinate $x \in [0, \pi/2)$. The dynamics of the system come from the Einstein and Klein-Gordon equations:

$$G_{\mu\nu} + \Lambda g_{\mu\nu} = 8\pi \left(\nabla_\mu \phi \nabla_\nu \phi - \frac{1}{2} g_{\mu\nu} (\nabla^\rho \phi \nabla_\rho \phi + m^2 \phi^2) \right), \quad \text{and} \quad \nabla^2 \phi - m^2 \phi = 0, \quad (2.2)$$

with Λ as cosmological constant for AdS, $\Lambda = -d(d-1)/2$.

Perturbing around static AdS, the scalar field is expanded in of odd powers of epsilon

$$\phi(t, x) = \epsilon \phi_1 + \epsilon^3 \phi_3 + \dots \quad (2.3)$$

and the metric functions A and δ in even powers,

$$A(t, x) = 1 + A_2 \epsilon^2 + \dots \quad \text{and} \quad \delta(t, x) = \epsilon^2 \delta_2 + \dots \quad (2.4)$$

We choose to work in the boundary gauge, where $\delta(t, \pi/2) = 0$, for reasons that we discuss below.

At linear order, ϕ_1 satisfies

$$\partial_t^2 \phi_1 + \hat{L} \phi_1 = 0 \quad \text{where} \quad \hat{L} \equiv \frac{1}{\mu} (\mu' \partial_x + \mu \partial_x^2) - \frac{m^2}{\cos^2(x)}, \quad (2.5)$$

where $\mu \equiv \tan^{d-1}(x)$. Writing the scalar field as the product of time- and position-dependent parts,

$$\phi_1(t, x) = \sum_j c_j(t) e_j(x), \quad (2.6)$$

we find that the basis functions $e_j(x)$ are the solutions to the eigenvalue equation

$$\hat{L}e_j(x) = \omega_j^2 e_j(x). \quad (2.7)$$

The general solution to this eigenvalue equation involves two types of functions: those that vanish as $x \rightarrow \pi/2$ and therefore are normalizable, and those that approach a finite value on the boundary and are non-normalizable. In many previous works, the dynamics of scalar fields have been studied using exclusively normalizable functions (as is the focus of § 3); however, we now wish to consider the effects of exciting the second class of basis functions. These driven systems have a time-dependent source term on the boundary that sends energy into bulk space through coupling with non-normalizable modes. Energy is then distributed to both normalizable and non-normalizable scalar modes, as well as the metric functions.

Writing out the full solution to (2.7) involves summing over both normalizable and non-normalizable eigenmodes [2],

$$\phi_1(t, x) = \sum_j c_j(t) e_j(x) + \sum_\alpha \bar{A}_\alpha(t) E_\alpha(x). \quad (2.8)$$

The values of \bar{A}_α are set by the choice of boundary conditions. For example, if the driving term on the boundary is a single, periodic function, then

$$\phi_1(t, \pi/2) = \mathcal{A} \cos \bar{\omega} t. \quad (2.9)$$

Extending the boundary condition to include the addition of two driving terms – or a more general sum over Fourier modes – would set further \bar{A}_α values.

The normalizable modes have eigenfunctions given by

$$e_j(x) = k_j (\cos(x))^{\Delta^+} P_j^{(d/2-1, \Delta^+-d/2)}(\cos(2x)) \quad (2.10)$$

$$k_j = 2 \sqrt{\frac{(j + \Delta^+/2)\Gamma(j+1)\Gamma(j + \Delta^+)}{\Gamma(j + d/2)\Gamma(j + \Delta^+ - d/2 + 1)}}, \quad (2.11)$$

with fully resonant eigenvalues $\omega_j = 2j + \Delta^+$ with $j \in \mathbb{Z}^*$. We define Δ^+ as the positive root of $\Delta(\Delta - d) = m^2$. This resonant spectrum is responsible for the secular growth of terms at third order, which must be controlled the secular term resummation procedure described below. On the other hand, the non-normalizable eigenfunctions have arbitrary frequencies ω_α , and are given by

$$E_\alpha(x) = (\cos(x))^{\Delta^+} {}_2F_1\left(\frac{\Delta^+ + \omega_\alpha}{2}, \frac{\Delta^+ - \omega_\alpha}{2}, d/2; \sin^2(x)\right). \quad (2.12)$$

By allowing frequencies and basis functions to remain unspecified for the time being, we can show that the $\mathcal{O}(\epsilon^3)$ part of the scalar field satisfies the equation

$$\ddot{\phi}_3 + \hat{L}\phi_3 = S = 2(A_2 - \delta_2)\ddot{\phi}_1 + (\dot{A}_2 - \dot{\delta}_2)\dot{\phi}_1 + (A'_2 - \delta'_2)\phi'_1 + m^2 A_2 \phi_1 \sec^2 x. \quad (2.13)$$

Following the steps outlined in Appendix A, we project (2.13) onto the basis of normalizable modes and, employing the solution $c_i(t) = a_i \cos(\omega_i t + b_i) = a_i \cos \theta_i$ for the time-dependent portion of the scalar field, find that the general expression for the source term is

$$\begin{aligned}
S_\ell = & \frac{1}{4} \sum_{\substack{i,j,k \\ k \neq \ell}}^{\infty} \frac{a_i a_j a_k \omega_k}{\omega_\ell^2 - \omega_k^2} \left[Z_{ijk\ell}^-(\omega_i + \omega_j - 2\omega_k) \cos(\theta_i + \theta_j - \theta_k) - Z_{ijk\ell}^-(\omega_i + \omega_j + 2\omega_k) \cos(\theta_i + \theta_j + \theta_k) - \right. \\
& \left. + Z_{ijk\ell}^+(\omega_i - \omega_j + 2\omega_k) \cos(\theta_i - \theta_j + \theta_k) - Z_{ijk\ell}^+(\omega_i - \omega_j - 2\omega_k) \cos(\theta_i - \theta_j - \theta_k) \right] \\
& + \frac{1}{2} \sum_{\substack{i,j,k \\ i \neq j}}^{\infty} a_i a_j a_k \omega_j \left(H_{ijk\ell} + m^2 V_{jki\ell} - 2\omega_k^2 X_{ijk\ell} \right) \left[\frac{1}{\omega_i - \omega_j} (\cos(\theta_i - \theta_j - \theta_k) + \cos(\theta_i - \theta_j + \theta_k)) \right. \\
& \left. - \frac{1}{\omega_i + \omega_j} (\cos(\theta_i + \theta_j - \theta_k) + \cos(\theta_i + \theta_j + \theta_k)) \right] \\
& - \frac{1}{4} \sum_{i,j,k}^{\infty} a_i a_j a_k \left[(2\omega_j \omega_k X_{ijk\ell} + m^2 V_{ijk\ell}) \cos(\theta_i + \theta_j - \theta_k) - (2\omega_j \omega_k X_{ijk\ell} - m^2 V_{ijk\ell}) \cos(\theta_i - \theta_j - \theta_k) \right. \\
& \left. + (2\omega_j \omega_k X_{ijk\ell} + m^2 V_{ijk\ell}) \cos(\theta_i - \theta_j + \theta_k) - (2\omega_j \omega_k X_{ijk\ell} - m^2 V_{ijk\ell}) \cos(\theta_i + \theta_j + \theta_k) \right] \\
& + \frac{1}{4} \sum_{i,j}^{\infty} a_i a_j a_\ell \omega_\ell \left[\tilde{Z}_{ij\ell}^-(\omega_i + \omega_j - 2\omega_\ell) \cos(\theta_i + \theta_j - \theta_\ell) - \tilde{Z}_{ij\ell}^-(\omega_i + \omega_j + 2\omega_\ell) \cos(\theta_i + \theta_j + \theta_\ell) \right. \\
& \left. + \tilde{Z}_{ij\ell}^+(\omega_i - \omega_j + 2\omega_\ell) \cos(\theta_i - \theta_j + \theta_\ell) - \tilde{Z}_{ij\ell}^+(\omega_i - \omega_j - 2\omega_\ell) \cos(\theta_i - \theta_j - \theta_\ell) \right] \\
& - \frac{1}{4} \sum_{i,j}^{\infty} a_i^2 a_j \left(H_{iij\ell} + m^2 V_{jii\ell} - 2\omega_j^2 X_{iij\ell} \right) [\cos(2\theta_i - \theta_j) + \cos(2\theta_i + \theta_j)] \\
& - \frac{1}{2} \sum_{i,j}^{\infty} a_i^2 a_j \left(H_{iij\ell} + m^2 V_{jii\ell} - 2\omega_j^2 X_{iij\ell} + 4\omega_i^2 \omega_j^2 P_{j\ell i} + 2\omega_i^2 (M_{j\ell i} + m^2 Q_{j\ell i}) \right) \cos \theta_j. \quad (2.14)
\end{aligned}$$

Note that sums and restrictions on indices must be interpreted as sums and restrictions on *frequencies* when any of the modes is non-normalizable, since $\omega_\alpha \neq 2\alpha + \Delta^+$ in general.

As mentioned above, the linear growth of resonant terms with time, i.e. secular growth, at $\mathcal{O}(\epsilon^3)$ can be absorbed into the time-dependent part of the scalar field at that order [3]. Thus, (2.13) tells us that

$$\ddot{c}_\ell^{(3)}(t) + \omega_\ell^2 c_\ell^{(3)}(t) = S_\ell^{(3)} \cos(\omega_\ell t + \varphi_\ell), \quad (2.15)$$

where $S_\ell^{(3)}$ is a polynomial in a_i determined by evaluating the resonant contributions from (2.14), and φ_ℓ is some combination of the b_i . To obtain the renormalization flow equations, we can rewrite the amplitudes and phases in terms of renormalized integration constants that exactly cancel the secular terms at each moment. Doing so yields the renormalization flow

equations for the renormalized constants [4]

$$\frac{2\omega_\ell}{\epsilon^2} \frac{da_\ell}{dt} = -S_\ell^{(3)} \sin(b_\ell - \varphi_\ell) , \quad (2.16)$$

$$\frac{2\omega_\ell a_\ell}{\epsilon^2} \frac{db_\ell}{dt} = -S_\ell^{(3)} \cos(b_\ell - \varphi_\ell) . \quad (2.17)$$

Note that the amplitudes and phases evolve with respect to the “slow time” $\tau = \epsilon^2 t$. In practice, once these flow equations can be written down the perturbative evolution of the system is determined up to a timescale of $t \sim \epsilon^{-2}$.

Let us now examine the exact form of $S_\ell^{(3)}$ for the case of a massive scalar field. As an exercise, we first derive the resonant contributions when the boundary source is zero, and therefore only normalizable modes are present. These results agree numerically with previous work on normalizable modes for massless scalars in the interior time gauge ($\delta(t, 0) = 0$) [5]. The definitions of the integral functions Z , H , X , etc. differ slightly from other works – in part because of the gauge choice, and in part because of a desire to separate out mass-dependent terms – and so are given explicitly in Appendix A.

3 Resonances From Normalizable Solutions

Consider the case where each of the basis functions are given by normalizable solutions. After time-averaging, the resonant contributions occur for the following combination of normalizable frequencies:

$$\omega_i \pm \omega_j \pm \omega_k = \pm \omega_\ell \quad (3.1)$$

which can be separated into the three distinct cases

$$\omega_i + \omega_j + \omega_k = \omega_\ell \quad (+ + +) \quad (3.2)$$

$$\omega_i - \omega_j - \omega_k = \omega_\ell \quad (+ - -) \quad (3.3)$$

$$\omega_i + \omega_j - \omega_k = \omega_\ell \quad (+ + -) . \quad (3.4)$$

We will see that the first two resonances, $(+ + +)$ and $(+ - -)$, will non-trivially vanish whenever their respective resonance conditions are satisfied. This is in agreement with the results shown for the massless scalar in the interior time gauge (as they must be, since the choice of time gauge should not change the existence of resonant channels). Here we include the expressions for the naturally vanishing resonances, choosing to explicitly express the mass dependence.

3.1 Naturally Vanishing Resonances: $(+ + +)$ and $(+ - -)$

Resonant contributions that come from the condition $\omega_i + \omega_j + \omega_k = \omega_\ell$ contribute to the total source term via

$$S_\ell = \underbrace{\sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \sum_{k=0}^{\infty}}_{\omega_i + \omega_j + \omega_k = \omega_\ell} \Omega_{ijk\ell} a_i a_j a_k \cos(\theta_i + \theta_j + \theta_k) + \dots , \quad (3.5)$$

where the dots denote other resonances. $\Omega_{ijk\ell}$ is given by

$$\begin{aligned}
\Omega_{ijk\ell} = & -\frac{1}{12}H_{ijk\ell}\frac{\omega_j(\omega_i+\omega_k+2\omega_j)}{(\omega_i+\omega_j)(\omega_j+\omega_k)} - \frac{1}{12}H_{ikj\ell}\frac{\omega_k(\omega_i+\omega_j+2\omega_k)}{(\omega_i+\omega_k)(\omega_j+\omega_k)} - \frac{1}{12}H_{jik\ell}\frac{\omega_i(\omega_j+\omega_k+2\omega_i)}{(\omega_i+\omega_j)(\omega_i+\omega_k)} \\
& - \frac{m^2}{12}V_{ijk\ell}\left(1+\frac{\omega_j}{\omega_j+\omega_k}+\frac{\omega_i}{\omega_i+\omega_k}\right) - \frac{m^2}{12}V_{jik\ell}\left(1+\frac{\omega_j}{\omega_i+\omega_j}+\frac{\omega_k}{\omega_i+\omega_k}\right) \\
& - \frac{m^2}{12}V_{kij\ell}\left(1+\frac{\omega_i}{\omega_i+\omega_j}+\frac{\omega_k}{\omega_j+\omega_k}\right) + \frac{1}{6}\omega_j\omega_kX_{ijk\ell}\left(1+\frac{\omega_j}{\omega_i+\omega_k}+\frac{\omega_k}{\omega_i+\omega_j}\right) \\
& + \frac{1}{6}\omega_i\omega_kX_{jik\ell}\left(1+\frac{\omega_i}{\omega_j+\omega_k}+\frac{\omega_k}{\omega_i+\omega_j}\right) + \frac{1}{6}\omega_i\omega_jX_{kij\ell}\left(1+\frac{\omega_i}{\omega_j+\omega_k}+\frac{\omega_j}{\omega_i+\omega_k}\right) \\
& - \frac{1}{12}Z_{ijk\ell}^-\left(\frac{\omega_k}{\omega_i+\omega_j}\right) - \frac{1}{12}Z_{ikj\ell}^-\left(\frac{\omega_j}{\omega_i+\omega_k}\right) - \frac{1}{12}Z_{jik\ell}^-\left(\frac{\omega_i}{\omega_j+\omega_k}\right). \tag{3.6}
\end{aligned}$$

The second naturally vanishing resonance comes from the condition $\omega_i - \omega_j - \omega_k = \omega_\ell$, and contributes to the total source term via

$$S_\ell = \sum_{j=0}^{\infty} \sum_{k=0}^{\infty} \Gamma_{(j+k+\ell+\Delta^+)jkl} a_j a_k a_{(j+k+\ell+\Delta^+)} \cos(\theta_{(j+k+\ell+\Delta^+)} - \theta_j - \theta_k) + \dots, \tag{3.7}$$

where

$$\begin{aligned}
\Gamma_{ijk\ell} = & \frac{1}{4}H_{ijk\ell}\frac{\omega_j(\omega_k-\omega_i+2\omega_j)}{(\omega_i-\omega_j)(\omega_j+\omega_k)} + \frac{1}{4}H_{jik\ell}\frac{\omega_k(\omega_j-\omega_i+2\omega_k)}{(\omega_i-\omega_k)(\omega_j+\omega_k)} + \frac{1}{4}H_{kij\ell}\frac{\omega_i(\omega_j+\omega_k-2\omega_i)}{(\omega_i-\omega_j)(\omega_i-\omega_k)} \\
& - \frac{1}{2}\omega_j\omega_kX_{ijk\ell}\left(\frac{\omega_k}{\omega_i-\omega_j}+\frac{\omega_j}{\omega_i-\omega_k}-1\right) + \frac{1}{2}\omega_i\omega_kX_{jik\ell}\left(\frac{\omega_k}{\omega_i-\omega_j}+\frac{\omega_i}{\omega_j+\omega_k}-1\right) \\
& + \frac{1}{2}\omega_i\omega_jX_{kij\ell}\left(\frac{\omega_j}{\omega_i-\omega_k}+\frac{\omega_i}{\omega_j+\omega_k}-1\right) + \frac{m^2}{4}V_{jik\ell}\left(\frac{\omega_j}{\omega_i-\omega_j}+\frac{\omega_k}{\omega_i-\omega_k}-1\right) \\
& - \frac{m^2}{4}V_{kij\ell}\left(\frac{\omega_i}{\omega_i-\omega_j}+\frac{\omega_k}{\omega_j+\omega_k}+1\right) - \frac{m^2}{4}V_{ijk\ell}\left(\frac{\omega_i}{\omega_i-\omega_k}+\frac{\omega_j}{\omega_j+\omega_k}+1\right) \\
& + \frac{1}{4}Z_{kij\ell}^-\left(\frac{\omega_i}{\omega_j+\omega_k}\right) - \frac{1}{4}Z_{ijk\ell}^+\left(\frac{\omega_k}{\omega_i-\omega_j}\right) - \frac{1}{4}Z_{jik\ell}^+\left(\frac{\omega_j}{\omega_i-\omega_k}\right). \tag{3.8}
\end{aligned}$$

Building on the work done with massless scalars, we are able to demonstrate that (3.6) and (3.8) continue to vanish for massive scalars ($m^2 \geq m_{BF}^2$) in the boundary gauge; thus, the dynamics governing the weakly turbulent transfer of energy is determined only from the remaining resonance channel. When non-normalizable modes are introduced, we will see that naturally vanishing resonances are not present and so the total third-order source term is the sum over all resonant channels.

3.2 $(++-)$

The first non-vanishing contributions arise when $\omega_i + \omega_j = \omega_k + \omega_\ell$. This contribution can be split into three coefficients that are evaluated for certain subsets of the allowed values for

the indices, namely

$$S_\ell = T_\ell a_\ell^3 \cos(\theta_\ell + \theta_\ell - \theta_\ell) + \sum_{i \neq \ell}^\infty R_{i\ell} a_i^2 a_\ell \cos(\theta_i + \theta_\ell - \theta_i) \\ + \sum_{i \neq \ell}^\infty \sum_{j \neq \ell}^\infty S_{ij(i+j-\ell)\ell} a_i a_j a_{(i+j-\ell)} \cos(\theta_i + \theta_j - \theta_{i+j-\ell}), \quad (3.9)$$

where each of the coefficients is given by

$$S_{ijk\ell} = -\frac{1}{4} H_{kij\ell} \frac{\omega_i(\omega_j - \omega_k + 2\omega_i)}{(\omega_i - \omega_k)(\omega_i + \omega_j)} - \frac{1}{4} H_{ij\ell k} \frac{\omega_j(\omega_i - \omega_k + 2\omega_j)}{(\omega_j - \omega_k)(\omega_i + \omega_j)} - \frac{1}{4} H_{jki\ell} \frac{\omega_k(\omega_i + \omega_j - 2\omega_k)}{(\omega_i - \omega_k)(\omega_j - \omega_k)} \\ - \frac{1}{2} \omega_j \omega_k X_{ijk\ell} \left(\frac{\omega_j}{\omega_i - \omega_k} - \frac{\omega_k}{\omega_i + \omega_j} + 1 \right) - \frac{1}{2} \omega_i \omega_k X_{jki\ell} \left(\frac{\omega_i}{\omega_j - \omega_k} - \frac{\omega_k}{\omega_i + \omega_j} + 1 \right) \\ + \frac{1}{2} \omega_i \omega_j X_{kij\ell} \left(\frac{\omega_i}{\omega_j - \omega_k} + \frac{\omega_j}{\omega_i - \omega_k} + 1 \right) - \frac{m^2}{4} V_{ijk\ell} \left(\frac{\omega_i}{\omega_i - \omega_k} + \frac{\omega_j}{\omega_j - \omega_k} + 1 \right) \\ + \frac{m^2}{4} V_{jki\ell} \left(\frac{\omega_k}{\omega_i - \omega_k} - \frac{\omega_j}{\omega_i + \omega_j} - 1 \right) + \frac{m^2}{4} V_{kij\ell} \left(\frac{\omega_k}{\omega_j - \omega_k} - \frac{\omega_i}{\omega_i + \omega_j} - 1 \right) \\ + \frac{1}{4} Z_{ijk\ell}^- \left(\frac{\omega_k}{\omega_i + \omega_j} \right) + \frac{1}{4} Z_{ikj\ell}^+ \left(\frac{\omega_j}{\omega_i - \omega_k} \right) + \frac{1}{4} Z_{jki\ell}^+ \left(\frac{\omega_i}{\omega_j - \omega_k} \right), \quad (3.10)$$

$$R_{i\ell} = \left(\frac{\omega_i^2}{\omega_\ell^2 - \omega_i^2} \right) (Y_{i\ell\ell i} - Y_{i\ell i\ell} + \omega_\ell^2 (X_{i\ell i\ell} - X_{\ell i i\ell})) + \left(\frac{\omega_i^2}{\omega_\ell^2 - \omega_i^2} \right) (H_{\ell i i\ell} + m^2 V_{i\ell\ell} - 2\omega_i^2 X_{\ell i i\ell}) \\ - \left(\frac{\omega_\ell^2}{\omega_\ell^2 - \omega_i^2} \right) (H_{i\ell i\ell} + m^2 V_{\ell i i\ell} - 2\omega_i^2 X_{i\ell i\ell}) - \frac{m^2}{4} (V_{i\ell i\ell} + V_{i\ell\ell}) + \omega_i^2 \omega_\ell^2 (P_{i\ell} - 2P_{\ell i}) \\ - \omega_i \omega_\ell X_{i\ell i\ell} - \frac{3m^2}{2} V_{\ell i i\ell} - \frac{1}{2} H_{i\ell\ell} + \omega_\ell^2 B_{i\ell} - \omega_i^2 M_{\ell i} - m^2 \omega_i^2 Q_{\ell i}, \quad (3.11)$$

and

$$T_\ell = \frac{1}{2} \omega_\ell^2 (X_{\ell\ell\ell\ell} + 4B_{\ell\ell} - 2M_{\ell\ell} - 2m^2 Q_{\ell\ell}) - \frac{3}{4} (H_{\ell\ell\ell\ell} + 3m^2 V_{\ell\ell\ell}). \quad (3.12)$$

Following the form of (2.16) - (2.17), these resonant terms set the evolution of the renormalized integration coefficients to be [6]

$$\frac{2\omega_\ell}{\epsilon^2} \frac{da_\ell}{dt} = - \sum_{i \neq \ell}^\infty \sum_{j \neq \ell}^\infty S_{ij(i+j-\ell)\ell} a_i a_j a_{(i+j-\ell)} \cos(b_\ell + b_{(i+j-\ell)} - b_i - b_j) \quad (3.13)$$

$$\frac{2\omega_\ell a_\ell}{\epsilon^2} \frac{db_\ell}{dt} = -T_\ell a_\ell^3 - \sum_{i \neq \ell}^\infty R_{i\ell} a_i^2 a_\ell \\ - \sum_{i \neq \ell}^\infty \sum_{j \neq \ell}^\infty S_{ij(i+j-\ell)\ell} a_i a_j a_{(i+j-\ell)} \cos(b_\ell + b_{(i+j-\ell)} - b_i - b_j) \quad (3.14)$$

To examine the effects of non-zero masses on R , S , and T – and therefore the differential equations that determine the evolution of the time-dependent coefficients $a_i(t)$ and $b_i(t)$ – we evaluate (3.10)-(3.12) for tachyonic, massless, and massive scalars in figure 1.

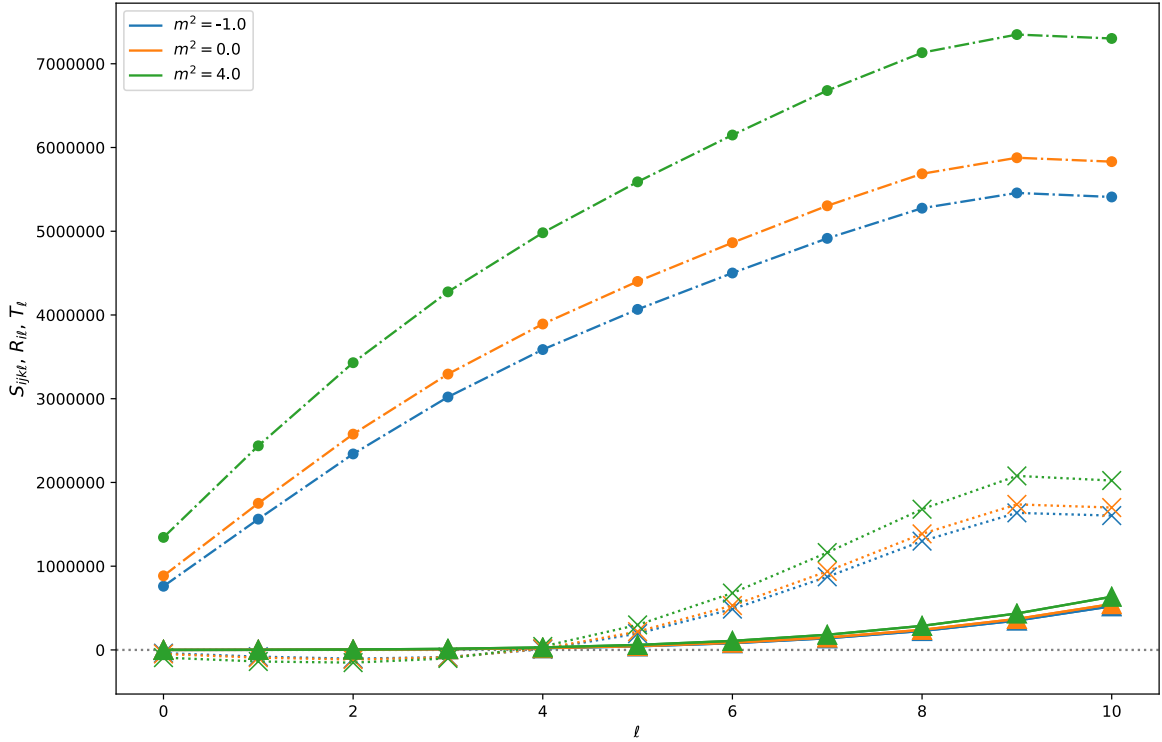


Figure 1: Evaluating (3.10)-(3.12) over different values of m^2 for $\ell \leq 10$. $S_{ij(i+j-\ell)\ell}$ is denoted by filled circles connected by dash-dotted lines; $R_{i\ell}$ is denoted by filled triangles connected by solid lines; T_ℓ is denoted by large Xs connected by dotted lines. Different values of m^2 are denoted by the colour of each series.

4 Resonances From Non-normalizable Modes

Now let us consider the excitation of non-normalizable modes by a driving term on the boundary of AdS.

Discuss falloff of A2 and delta2 with three NN modes, but don't calculate anything new. Mention overlap of NN and normalizable cases when $\omega(i) \pm \omega(j) = \text{NN frequency}$. Then focus on two NN modes.

Take the boundary condition to be at $\mathcal{O}(\epsilon)$, then at 3rd order we must have $\phi_3(x \rightarrow \pi/2) = 0$, which sets ω_ℓ to be normalizable. One could take all contributions to the 3rd order source term to be normalizable, but then the result is the same as the previous section but with a different first-order expression. When one mode is non-normalizable, the third-order boundary condition $\phi_3(t, \pi/2) = 0$ is violated; thus, at least two NN modes must be present so that the boundary condition can be met. When three non-normalizable modes exist, there are two possibilities: first, that any combination of generically non-integer frequencies gives a non-integer value and so vanishes when projected onto ω_ℓ and time averaged; second, some

particular combination of the non-normalizable frequencies gives an integer frequency, in which case the procedure for determining the contribution to S_ℓ follows the same procedure as the all-normalizable case. Therefore, the pertinent question is what contribution to the source term results from two of $\{\omega_i, \omega_j, \omega_k\}$ being non-normalizable. Because this choice breaks some of the symmetries that contributed to the previous expressions for resonance channels, the resonance conditions must be re-examined starting from the source expression (2.14).

Before proceeding further, an important consideration what the effect of non-normalizable modes are on the perturbative expansion that leads to the source equations. Since non-normalizable solutions do not have well-defined norms, we do not know *a priori* that the inner products described in Appendix A are still finite. To investigate this, consider the generic expression for the second-order metric function

$$A_2 = -\nu \int_0^x dy \mu \left((\dot{\phi}_1)^2 + (\phi_1')^2 + m^2 \phi_1^2 \sec^2 x \right), \quad (4.1)$$

in the limit of $x \rightarrow \pi/2$, and let the scalar field ϕ_1 be given by a generic superposition of normalizable and non-normalizable eigenfunctions as in (2.8). Ignoring the time-dependent contributions, we find that

$$\lim_{\tilde{x} \rightarrow 0} A_2(\tilde{x} \equiv \pi/2 - x) = \tilde{x}^{-\xi} \left(\frac{2\tilde{x}^{2+d}}{2-\xi} - \frac{\tilde{x}^d(1+(\Delta^-)^2)}{\xi} \right) \quad (4.2)$$

where we have defined $\xi = \sqrt{d^2 + 4m^2}$. In the massless case, $\xi = d$ and all powers of \tilde{x} are non-negative; thus, the limit is finite. For tachyonic masses of $m_{BF}^2 < m^2 < 0$, $0 < \xi < d$ and the limit is again finite. However, for scalars that either saturate the Breitenlohner-Freedman bound, or have $m^2 > 0$, part of the limit diverges. In order for the boundary to remain asymptotically AdS, counter-terms in the bulk action would be required to cancel such divergences – a case we will not address presently. Thus, we will restrict our discussion to $m_{BF}^2 < m^2 \leq 0$ to avoid these issues. A similar check on the near-boundary behaviour of δ_2 shows that the gauge condition $\delta_2(t, x = \pi/2)$ remains unchanged by the addition of non-normalizable modes given the same restrictions on the mass of the scalar field. With these restrictions in mind, let us now examine the resonances produced by the activation of non-normalizable modes.

4.1 Two Non-normalizable Modes with Equal Frequencies

As a first case, let us assume that the two non-normalizable modes have equal, constant, and arbitrary frequencies, $\bar{\omega}$. Applying the time-averaging procedure to the source S_ℓ once again eliminates all contributions except those that satisfy (3.1). We are now free to choose any one of $\{\omega_i, \omega_j, \omega_k\}$ to be normalizable and consider when the resonance condition is satisfied.

In particular, we find that the following combinations are resonant:

$$\omega_i - \omega_j + \omega_k - \omega_\ell = 0 \quad \Rightarrow \quad \text{either } \omega_i \text{ or } \omega_k \text{ is normalizable} \quad (4.3)$$

$$\omega_i + \omega_j - \omega_k - \omega_\ell = 0 \quad \Rightarrow \quad \text{either } \omega_i \text{ or } \omega_j \text{ is normalizable} \quad (4.4)$$

$$\omega_i - \omega_j - \omega_k + \omega_\ell = 0 \quad \Rightarrow \quad \text{either } \omega_j \text{ or } \omega_k \text{ is normalizable.} \quad (4.5)$$

When any of these resonance conditions is met, the remaining normalizable mode will have a frequency equal to ω_ℓ , collapsing all sums over frequencies so that

$$S_\ell = \bar{T}_\ell a_\ell \bar{A}_\omega^2 \cos(\theta_\ell) + \dots, \quad (4.6)$$

where the non-normalizable modes their amplitudes \bar{A}_ω set by the choice of boundary condition. Collecting the appropriate terms in (2.14), and evaluating the each possible resonance we find that

$$\begin{aligned} \bar{T}_\ell = & \left[\frac{1}{2} Z_{\ell\omega\omega}^- \left(\frac{\bar{\omega}}{\omega_\ell + \bar{\omega}} \right) + \frac{1}{2} Z_{\ell\omega\omega}^+ \left(\frac{\bar{\omega}}{\omega_\ell - \bar{\omega}} \right) + H_{\ell\omega\omega} \left(\frac{\bar{\omega}^2}{\omega_\ell^2 - \bar{\omega}^2} \right) - H_{\ell\omega\omega} \left(\frac{\omega_\ell^2}{\omega_\ell^2 - \bar{\omega}^2} \right) \right. \\ & - m^2 V_{\ell\omega\omega} \left(\frac{\omega_\ell^2}{\omega_\ell^2 - \bar{\omega}^2} \right) + m^2 V_{\omega\omega\ell} \left(\frac{\bar{\omega}^2}{\omega_\ell^2 - \bar{\omega}^2} \right) + 2X_{\omega\omega\ell} \left(\frac{\bar{\omega}^2 \omega_\ell^2}{\omega_\ell^2 - \bar{\omega}^2} \right) - 2X_{\ell\omega\omega} \left(\frac{\bar{\omega}^4}{\omega_\ell^2 - \bar{\omega}^2} \right) \Big]_{\bar{\omega} \neq \omega_\ell} \\ & + \omega_\ell^2 X_{\omega\omega\ell} - \bar{\omega}^2 X_{\ell\omega\omega} - \frac{3}{2} m^2 V_{\ell\omega\omega} - \frac{1}{2} m^2 V_{\omega\omega\ell} - \frac{1}{2} H_{\omega\omega\ell} + \omega_\ell^2 \tilde{Z}_{\omega\omega\ell}^+ - 2\bar{\omega}^2 \omega_\ell^2 P_{\ell\omega\omega} \\ & - \bar{\omega}^2 (\omega_\ell^2 P_{\ell\omega\omega} - B_{\ell\omega\omega}) . \end{aligned} \quad (4.7)$$

Notice that the terms in the square braces only contribute when $\bar{\omega} \neq \omega_\ell$. Beginning from (2.14), only terms in the square braces that are proportional to Z^\pm are limited in this way; the remaining terms have no such restriction. However, it can be shown that integral functions with permuted indices are equal when the non-normalizable frequency equals the normalizable frequency. Upon simplification, factors of $\omega_\ell^2 - \bar{\omega}^2$ are canceled and the overall contribution to T_ℓ from the terms in the braces is zero. Thus, these terms are grouped with those that have natural restrictions on the indices.

In figures 2 and 3, we evaluate (4.7) for $\ell < 10$ over a variety of $\bar{\omega}$ values first for a massless scalar, then for a tachyonic scalar.

Other resonant contributions become possible for more restrictive values of the non-normalizable frequency, such as if $\bar{\omega}$ is allowed to be an integer. These contributions are denoted by the dots in (4.6) and are discussed briefly in Appendix B.

4.2 Special Values of Non-normalizable Frequencies

Let us now consider special values of non-normalizable frequencies that will lead to a greater number of resonance channels. While general non-normalizable frequencies do not require any such restrictions, we will find it informative to examine these special cases as they possess more symmetry in index/frequency values than the case of equal non-normalizable frequencies, but less than all-normalizable modes.

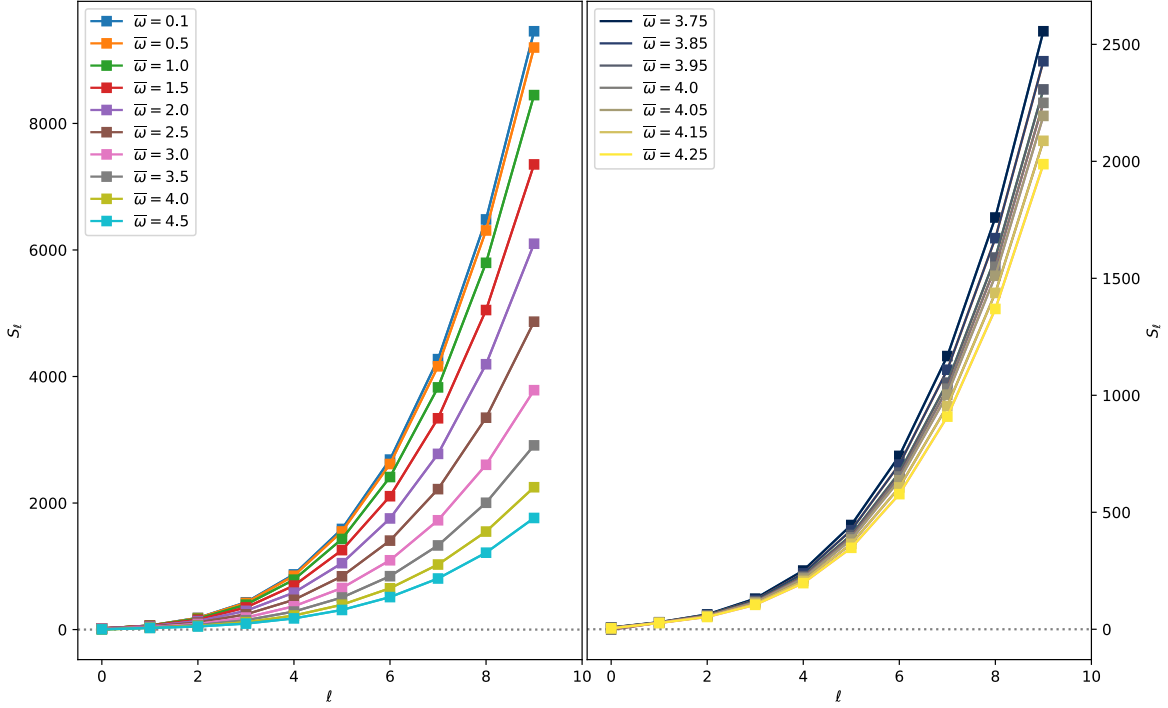


Figure 2: *Left:* Evaluating (4.7) when $m^2 = 0$ for various choices of $\bar{\omega}$. *Right:* The behaviour of S_ℓ for $\bar{\omega}$ values near ω_0 .

4.2.1 Add to an integer

First, we choose two of the modes to be non-normalizable with frequencies $\bar{\omega}_1$ and $\bar{\omega}_2$ that add to give an integer: $\bar{\omega}_1 + \bar{\omega}_2 = 2n$ where $n = 1, 2, 3, \dots$ (note that the $n = 0$ case means that both $\bar{\omega}_1$ and $\bar{\omega}_2$ would need to be zero by the positive-frequency requirement and so would not contribute). Furthermore, either frequency need not be an integer and therefore the difference $|\bar{\omega}_1 - \bar{\omega}_2|$ will, in general, not be an integer. In §4.3, we examine the case when the difference of non-normalizable frequencies is an integer.

When we consider possible resonance channels, we see that resonances can be grouped into

$$(++) : \omega_i + 2n = \omega_\ell \quad \forall \ell \geq n \quad (4.8)$$

$$(+ -) : \omega_i - 2n = \omega_\ell \quad \forall n \quad (4.9)$$

for any $m_{BF}^2 < m^2 < 0$. However, for a massless scalar, we have an additional channel

$$(-+) : -\omega_i + 2n = \omega_\ell \quad \forall n \geq \ell + d \quad (4.10)$$

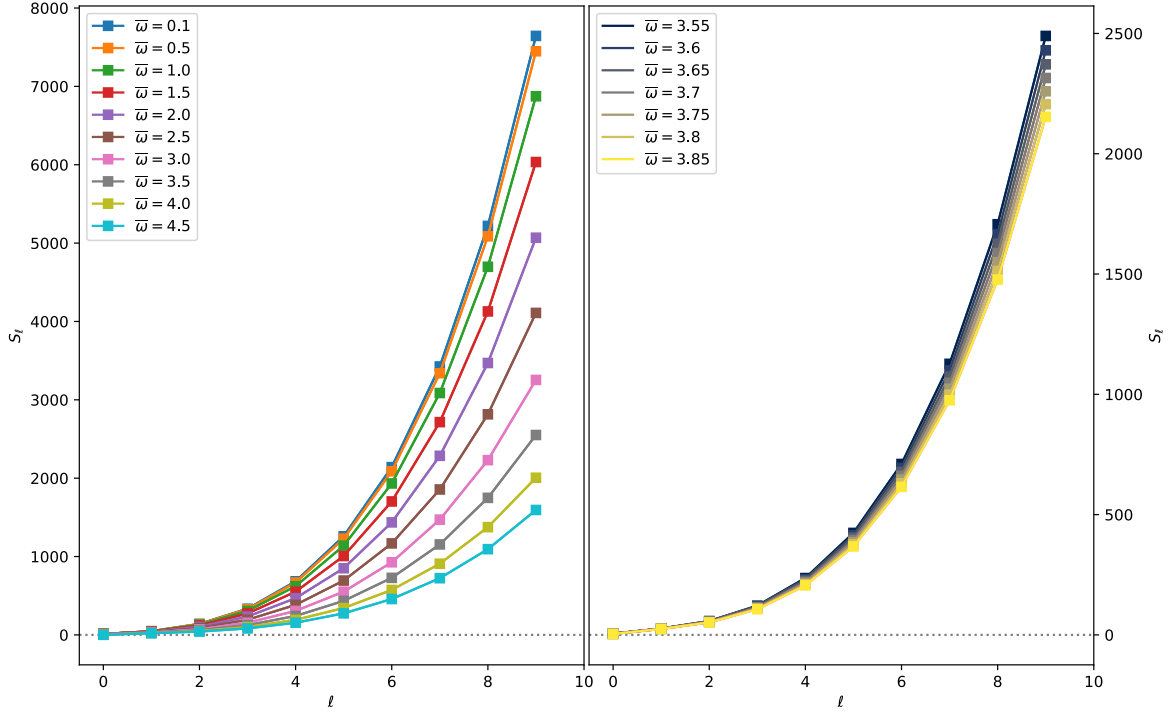


Figure 3: *Left:* Evaluating \bar{T}_ℓ for a tachyon with $m^2 = -1.0$. *Right:* The behaviour of S_ℓ near $\omega_0 = \Delta^+ \approx 3.7$.

Adding the channels together, the total source term is

$$\begin{aligned}
S_\ell = & \bar{R}_{(\ell+n)\ell}^{(+)} \bar{A}_1 \bar{A}_2 a_{(\ell+n)} \cos(\theta_{(\ell+n)} - 2nt) + \bar{T}_\ell \bar{A}_1 \bar{A}_2 a_\ell \cos(\theta_\ell) \\
& + \left[\Theta(n - \ell - d) \bar{R}_{(n-\ell-d)\ell}^{(-)} \bar{A}_1 \bar{A}_2 a_{(n-\ell-d)} \cos(\theta_{(n-\ell-d)} - 2nt) \right]_{m^2=0} \\
& + \Theta(\ell - n) \bar{R}_{(\ell-n)\ell}^{(+)} \bar{A}_1 \bar{A}_2 a_{(\ell-n)} \cos(\theta_{(\ell-n)} + 2nt), \tag{4.11}
\end{aligned}$$

where the Heaviside step function $\Theta(x)$ enforces the restrictions on the indices in (4.8) and (4.10). In the preceding expressions, the sum over all $\bar{\omega}_1, \bar{\omega}_2$ such that $\bar{\omega}_1 + \bar{\omega}_2 = 2n$ is implied, and only the restrictions on individual frequencies are included. Examining each channel in

(4.11) individually, we find

$$\begin{aligned}
\overline{R}_{i\ell}^{(++)} = & -\frac{1}{4} \sum_{\overline{\omega}_2 \neq \omega_\ell} \frac{\overline{\omega}_2}{\omega_\ell - \overline{\omega}_2} Z_{i12\ell}^- - \frac{1}{4} \sum_{\overline{\omega}_1 \neq \omega_\ell} \frac{\overline{\omega}_1}{\omega_\ell - \overline{\omega}_1} Z_{i21\ell}^- - \frac{1}{8n} \sum (\omega_\ell - 2n) Z_{12i\ell}^- \\
& - \frac{1}{4} \sum_{\omega_i \neq \overline{\omega}_1} \frac{1}{\omega_\ell - \overline{\omega}_2} \left[\overline{\omega}_1 (H_{i12\ell} + m^2 V_{12i\ell} - 2\overline{\omega}_2^2 X_{i12\ell}) + (\omega_\ell - 2n) (H_{1i2\ell} + m^2 V_{i21\ell} - 2\overline{\omega}_2^2 X_{1i2\ell}) \right] \\
& - \frac{1}{4} \sum_{\omega_i \neq \overline{\omega}_2} \frac{1}{\omega_\ell - \overline{\omega}_1} \left[\overline{\omega}_2 (H_{i21\ell} + m^2 V_{21i\ell} - 2\overline{\omega}_1^2 X_{i21\ell}) + (\omega_\ell - 2n) (H_{2i1\ell} + m^2 V_{i12\ell} - 2\overline{\omega}_1^2 X_{2i1\ell}) \right] \\
& - \frac{1}{8n} \sum_{\overline{\omega}_1 \neq \overline{\omega}_2} \left[\overline{\omega}_1 H_{21i\ell} + \overline{\omega}_2 H_{12i\ell} + m^2 (\overline{\omega}_1 V_{1i2\ell} + \overline{\omega}_2 V_{2i1\ell}) - (\omega_\ell - 2n)^2 (\overline{\omega}_1 X_{21i\ell} + \overline{\omega}_2 X_{12i\ell}) \right] \\
& + \frac{1}{2} \sum \left[\overline{\omega}_1 \overline{\omega}_2 X_{i12\ell} + (\omega_\ell - 2n) (\overline{\omega}_1 X_{21i\ell} + \overline{\omega}_2 X_{12i\ell}) - \frac{m^2}{2} (V_{i12\ell} + V_{i21\ell} + V_{12i\ell}) \right]
\end{aligned} \tag{4.12}$$

The notation $X_{i12\ell}$ corresponds to evaluating X_{ijkl} with $\omega_j = \overline{\omega}_1$ and $\omega_k = \overline{\omega}_2$. Next, we find that

$$\begin{aligned}
\overline{R}_{i\ell}^{(+-)} = & -\frac{1}{4} \sum \left[\frac{(\omega_\ell + 2n)}{2n} Z_{12i\ell}^- + 2(\omega_\ell + 2n) (\overline{\omega}_1 X_{21i\ell} + \overline{\omega}_2 X_{12i\ell}) \right. \\
& - \frac{\overline{\omega}_1}{(\omega_\ell + \overline{\omega}_2)} (H_{i12\ell} + m^2 V_{12i\ell} - 2\overline{\omega}_2^2 X_{i12\ell}) + \frac{(\omega_\ell + 2n)}{(\omega_\ell + \overline{\omega}_2)} (H_{1i2\ell} + m^2 V_{i21\ell} - 2\overline{\omega}_2^2 X_{1i2\ell}) \\
& - \frac{\overline{\omega}_2}{(\omega_\ell + \overline{\omega}_1)} (H_{i21\ell} + m^2 V_{21i\ell} - 2\overline{\omega}_1^2 X_{i21\ell}) + \frac{(\omega_\ell + 2n)}{(\omega_\ell + \overline{\omega}_1)} (H_{2i1\ell} + m^2 V_{i12\ell} - 2\overline{\omega}_1^2 X_{2i1\ell}) \\
& \left. - 2\overline{\omega}_1 \overline{\omega}_2 X_{i12\ell} + m^2 (V_{12i\ell} + V_{i12\ell} + V_{i21\ell}) \right] + \frac{1}{4} \sum_{\overline{\omega}_2 \neq \omega_\ell} \frac{\overline{\omega}_1 \overline{\omega}_2 (\omega_\ell + 2n)}{\omega_\ell + \overline{\omega}_2} (X_{21i\ell} - X_{\ell i12}) \\
& + \frac{1}{4} \sum_{\overline{\omega}_1 \neq \omega_\ell} \frac{\overline{\omega}_1 \overline{\omega}_2 (\omega_\ell + 2n)}{\omega_\ell + \overline{\omega}_1} (X_{12i\ell} - X_{\ell i12}).
\end{aligned} \tag{4.13}$$

When $m^2 = 0$, we have contributions from

$$\begin{aligned}
\overline{R}_{i\ell}^{(-+)} = & \frac{1}{4} \sum_{\overline{\omega}_2 \neq \omega_\ell} \frac{\overline{\omega}_2}{\omega_\ell - \overline{\omega}_2} Z_{i12\ell}^+ + \frac{1}{4} \sum_{\overline{\omega}_1 \neq \omega_\ell} \frac{\overline{\omega}_1}{\omega_\ell - \overline{\omega}_1} Z_{i21\ell}^+ + \frac{1}{4} \sum_{i \neq \ell} \left(\frac{2n - \omega_\ell}{2n} \right) Z_{12i\ell}^- \\
& + \frac{1}{4} \sum_{\overline{\omega}_1 \neq \omega_i} \frac{1}{\omega_i - \overline{\omega}_1} \left[\overline{\omega}_1 (H_{i12\ell} - 2\overline{\omega}_2^2 X_{i12\ell}) - (2n - \omega_\ell) (H_{1i2\ell} - 2\overline{\omega}_2^2 X_{1i2\ell}) \right] \\
& + \frac{1}{4} \sum_{\overline{\omega}_2 \neq \omega_i} \frac{1}{\omega_i - \overline{\omega}_2} \left[\overline{\omega}_2 (H_{i21\ell} - 2\overline{\omega}_1^2 X_{i21\ell}) - (2n - \omega_\ell) (H_{2i1\ell} - 2\overline{\omega}_1^2 X_{2i1\ell}) \right] \\
& - \frac{1}{8n} \sum_{\overline{\omega}_1 \neq \overline{\omega}_2} \left[\overline{\omega}_1 H_{21i\ell} + \overline{\omega}_2 H_{12i\ell} - 2(2n - \omega_\ell)^2 (\overline{\omega}_1 X_{21i\ell} + \overline{\omega}_2 X_{12i\ell}) \right] \\
& - \frac{1}{2} \sum \left[(2n - \omega_\ell) (\overline{\omega}_1 X_{21i\ell} + \overline{\omega}_2 X_{12i\ell}) - \overline{\omega}_1 \overline{\omega}_2 X_{i12\ell} \right].
\end{aligned} \tag{4.14}$$

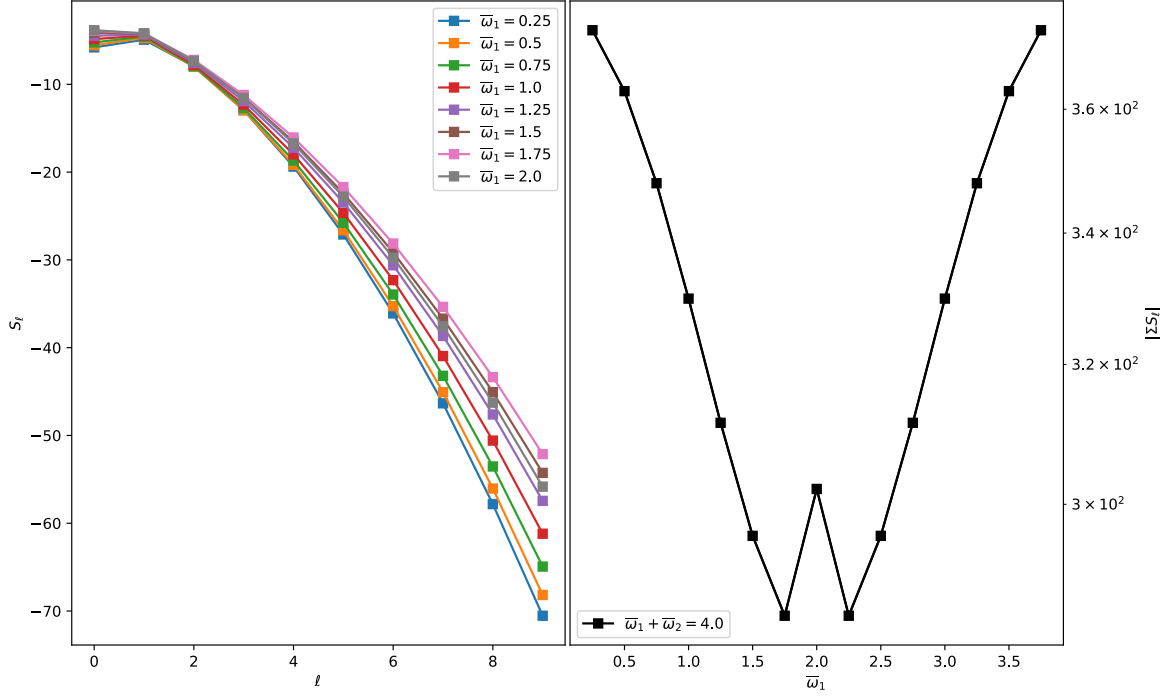


Figure 4: *Left:* Source term values for a tachyonic scalar with $m^2 = -1.0$ when the frequencies of non-normalizable modes sum to 4.0. *Right:* The absolute value of the sum of the source terms for each choice of $\bar{\omega}_1, \bar{\omega}_2$.

NB. In (4.14) only, $\omega_i = 2i + \Delta^+ = 2i + d$ since this term requires that $m^2 = 0$ to contribute. We maintain the same notation out of convenience, despite the special case.

Finally,

$$\begin{aligned} \bar{T}_\ell = & \frac{1}{2}\omega_\ell^2 \left(\tilde{Z}_{11\ell}^+ + \tilde{Z}_{22\ell}^+ \right) - \frac{1}{2} \left[H_{11\ell\ell} + H_{22\ell\ell} + m^2 (V_{\ell 11\ell} + V_{\ell 22\ell}) - 2\omega_\ell^2 (X_{11\ell\ell} + X_{22\ell\ell}) \right. \\ & \left. + 4\omega_\ell^2 (\bar{\omega}_1^2 P_{\ell\ell 1} + \bar{\omega}_2^2 P_{\ell\ell 2}) + 2\bar{\omega}_1^2 M_{\ell\ell 1} + 2\bar{\omega}_2^2 M_{\ell\ell 2} + 2m^2 (\bar{\omega}_1^2 Q_{\ell\ell 1} + \bar{\omega}_2^2 Q_{\ell\ell 2}) \right]. \end{aligned} \quad (4.15)$$

In figure 4, we compute the total source term – modulo the amplitudes a_i and \bar{A}_α – for a tachyonic scalar with $n = 2$. Figure 5 provides a comparison between the value of the source term for a massless scalar between two choices of n : one that includes contributions from $\bar{R}_{i\ell}^{(-+)}$ and one that does not.

4.3 Integer Plus χ

Finally, let us consider the case where the non-normalizable frequencies are non-integer, but differ from integer values by a set amount. In analogue to the case where all modes are normalizable, we consider the non-normalizable frequencies to be shifted away from integer

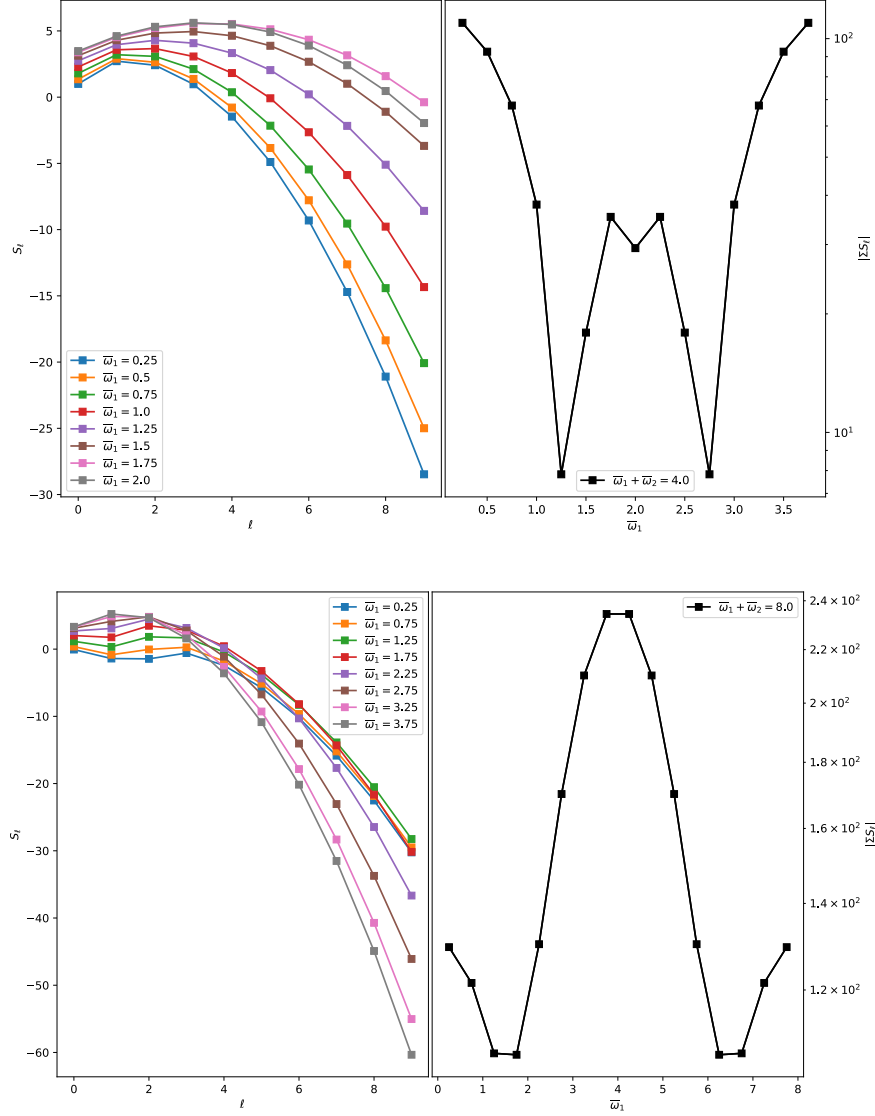


Figure 5: *Above:* The value of (4.11) as a function of ℓ for a massless scalar with values of $\bar{\omega}_1$ and $\bar{\omega}_2$ chosen so that $\bar{\omega}_1 + \bar{\omega}_2 = 4$. *Below:* The same plot but with values chosen to satisfy $\bar{\omega}_1 + \bar{\omega}_2 = 8$.

values by

$$\omega_\gamma = 2\gamma + \chi, \quad (4.16)$$

where $\gamma \in \mathbb{Z}^*$ (greek letters are chosen to differentiate these non-normalizable modes from normalizable modes with integer frequencies, which use roman letters). We furthermore limit χ to be non-integer¹ and set $m^2 = 0$ throughout. For this choice of non-normalizable

¹Indeed, for integer values of χ , the sum or difference of two non-normalizable modes could be an integer. This would either be covered by the work in § 4.2.1, or be a slight variation of it.

frequencies there are no resonant contributions from the all-plus channel, unlike the naturally vanishing resonance found in §3.1. Only when either $\omega_i + \omega_\gamma = \omega_\beta - \omega_\ell$, or $\omega_i + \omega_\gamma = \omega_\beta + \omega_\ell$ with $i + \gamma \geq \ell$, are resonant terms present. Let us examine each case separately.

4.3.1 $\omega_i + \omega_\gamma = \omega_\beta - \omega_\ell$

When the resonance condition $\omega_i + \omega_\gamma = \omega_\beta - \omega_\ell$ is met, the contribution to the source term is of the form

$$S_\ell = \sum_{i \neq \ell} \sum_{\gamma \neq \beta} \bar{S}_{i(i+\gamma+\ell)\gamma\ell} a_i \bar{A}_{(i+\gamma+\ell)} \bar{A}_\gamma \cos(\theta_i - \theta_{(i+\gamma+\ell)} + \theta_\gamma) \\ + \sum_{\beta} \bar{R}_{\beta\ell} a_\ell \bar{A}_\beta^2 \cos(\theta_\ell + \theta_\beta - \theta_\beta) + \dots, \quad (4.17)$$

where

$$\bar{S}_{i\beta\gamma\ell} = \frac{1}{4} H_{\beta\gamma i\ell} \frac{\omega_\gamma(\omega_i - \omega_\beta + 2\omega_\gamma)}{(\omega_\beta - \omega_\gamma)(\omega_i + \omega_\gamma)} - \frac{1}{4} H_{\gamma\beta i\ell} \frac{\omega_\beta(\omega_i + \omega_\gamma - 2\omega_\beta)}{(\omega_i - \omega_\beta)(\omega_\beta - \omega_\gamma)} - \frac{1}{4} H_{\gamma i\beta\ell} \frac{\omega_i(\omega_\gamma - \omega_\beta + 2\omega_i)}{(\omega_i - \omega_\beta)(\omega_i + \omega_\gamma)} \\ + \frac{1}{2} \omega_i \omega_\gamma X_{\beta\gamma i\ell} \left(\frac{\omega_\gamma}{\omega_i - \omega_\beta} - \frac{\omega_i}{\omega_\beta + \omega_\gamma} + 1 \right) + \frac{1}{2} \omega_i \omega_\beta X_{\gamma\beta i\ell} \left(\frac{\omega_i}{\omega_\beta - \omega_\gamma} + \frac{\omega_\beta}{\omega_i + \omega_\gamma} - 1 \right) \\ + \frac{1}{2} \omega_\beta \omega_\gamma X_{i\beta\gamma\ell} \left(\frac{\omega_\beta}{\omega_i + \omega_\gamma} - \frac{\omega_\gamma}{\omega_i - \omega_\beta} - 1 \right) - \frac{1}{4} Z_{\beta\gamma i\ell}^+ \left(\frac{\omega_i}{\omega_i + \omega_\ell} \right) \\ + \frac{1}{4} Z_{i\gamma\beta\ell}^- \left(\frac{\omega_\beta}{\omega_\ell - \omega_\beta} \right) + \frac{1}{4} Z_{i\beta\gamma\ell}^+ \left(\frac{\omega_\gamma}{\omega_\ell + \omega_\gamma} \right), \quad (4.18)$$

and

$$\bar{R}_{\beta\ell} = \frac{1}{4} Z_{\ell\beta\beta\ell}^- \left(\frac{\omega_\beta}{\omega_\ell + \omega_\beta} \right) + \frac{1}{4} Z_{\ell\beta\beta\ell}^+ \left(\frac{\omega_\beta}{\omega_\ell - \omega_\beta} \right) + \frac{1}{2} H_{\ell\beta\beta\ell} \left(\frac{\omega_\beta^2}{\omega_\ell^2 - \omega_\beta^2} \right) - \frac{1}{2} H_{\beta\ell\beta\ell} \left(\frac{\omega_\ell^2}{\omega_\ell^2 - \omega_\beta^2} \right) \\ + X_{\beta\ell\beta\ell} \left(\frac{\omega_\ell^4}{\omega_\ell^2 - \omega_\beta^2} \right) - \frac{1}{2} \omega_\beta^2 X_{\ell\beta\beta\ell} \left(\frac{\omega_\ell^2 + \omega_\beta^2}{\omega_\ell^2 - \omega_\beta^2} \right) - \frac{1}{2} H_{\ell\beta\beta\ell} + \omega_\ell^2 \tilde{Z}_{\beta\beta\ell}^+ - 2\omega_\beta^2 \omega_\ell^2 P_{\ell\ell\beta} - \omega_\beta^2 M_{\ell\ell\beta}. \quad (4.19)$$

4.3.2 $\omega_i + \omega_\gamma = \omega_\beta + \omega_\ell$

Similarly, when the resonance condition $\omega_i + \omega_\gamma = \omega_\beta + \omega_\ell$ is met, the contribution to the source term is

$$S_\ell = \underbrace{\sum_{i \neq \ell} \sum_{\gamma \neq \beta}}_{i+\gamma \geq \ell} \bar{S}_{i(i+\gamma-\ell)\gamma\ell} a_i \bar{A}_{(i+\gamma-\ell)} \bar{A}_\gamma \cos(\theta_i - \theta_{(i+\gamma-\ell)} + \theta_\gamma) \\ + \sum_{\beta} \bar{R}_{\beta\ell} a_\ell \bar{A}_\beta^2 \cos(\theta_\ell + \theta_\beta - \theta_\beta) + \dots, \quad (4.20)$$

where

$$\begin{aligned}
\bar{S}_{i\beta\gamma\ell} = & \frac{1}{4}H_{\beta\gamma i\ell} \frac{\omega_\gamma(\omega_i - \omega_\beta)}{(\omega_\beta - \omega_\gamma)(\omega_i - \omega_\gamma)} - \frac{1}{4}H_{\gamma\beta i\ell} \frac{\omega_\beta(\omega_\ell - \omega_\beta)}{(\omega_\beta - \omega_\gamma)(\omega_i - \omega_\beta)} + \frac{1}{4}H_{\beta i\gamma\ell} \frac{\omega_i(\omega_\gamma - \omega_\beta)}{(\omega_i - \omega_\beta)(\omega_i - \omega_\gamma)} \\
& + \frac{1}{2}\omega_i\omega_\gamma X_{\beta\gamma i\ell} \left(\frac{\omega_\gamma}{\omega_i - \omega_\beta} - \frac{\omega_i}{\omega_\beta - \omega_\gamma} + 1 \right) + \frac{1}{2}\omega_i\omega_\beta X_{\gamma\beta i\ell} \left(\frac{\omega_i}{\omega_\beta - \omega_\gamma} - \frac{\omega_\beta}{\omega_i - \omega_\gamma} - 1 \right) \\
& + \frac{1}{2}\omega_\beta\omega_\gamma X_{i\beta\gamma\ell} \left(\frac{\omega_\beta}{\omega_i - \omega_\gamma} - \frac{\omega_\gamma}{\omega_i - \omega_\beta} - 1 \right) + \frac{1}{4}Z_{i\gamma\beta\ell}^- \left(\frac{\omega_\beta}{\omega_\ell + \omega_\beta} \right) \\
& + \frac{1}{4}Z_{i\beta\gamma\ell}^+ \left(\frac{\omega_\gamma}{\omega_\ell - \omega_\gamma} \right) - \frac{1}{4}Z_{\beta\gamma i\ell}^+ \left(\frac{\omega_i}{\omega_i - \omega_\ell} \right), \tag{4.21}
\end{aligned}$$

and

$$\begin{aligned}
\bar{R}_{\beta\ell} = & \frac{1}{4}Z_{\ell\beta\beta\ell}^- \left(\frac{\omega_\beta}{\omega_\ell + \omega_\beta} \right) + \frac{1}{4}Z_{\ell\beta\beta\ell}^+ \left(\frac{\omega_\beta}{\omega_\ell - \omega_\beta} \right) + \frac{1}{2}H_{\ell\beta\beta\ell} \left(\frac{\omega_\beta^2}{\omega_\ell^2 - \omega_\beta^2} \right) - \frac{1}{2}H_{\beta\ell\beta\ell} \left(\frac{\omega_\ell^2}{\omega_\ell^2 - \omega_\beta^2} \right) \\
& + X_{\beta\beta\ell\ell} \left(\frac{\omega_\ell^2}{\omega_\ell^2 - \omega_\beta^2} \right) + \frac{1}{2}\omega_\beta^2 X_{\ell\beta\beta\ell} \left(\frac{\omega_\ell^2 + \omega_\beta^2}{\omega_\ell^2 - \omega_\beta^2} \right) - \frac{1}{2}H_{\beta\beta\ell\ell} + \omega_\ell^2 \tilde{Z}_{\beta\beta\ell}^+ - 2\omega_\beta^2\omega_\ell^2 P_{\ell\ell\beta} - \omega_\beta^2 M_{\ell\ell\beta} \tag{4.22}
\end{aligned}$$

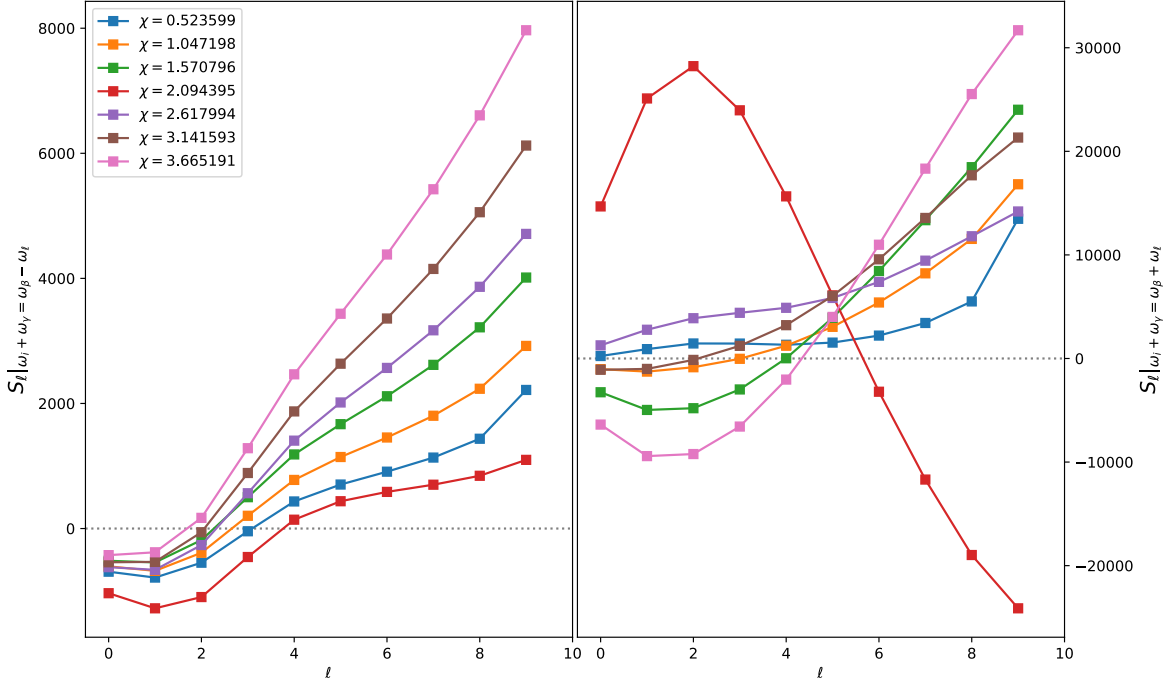


Figure 6: *Left:* Evaluating the source term (4.17) for various values of χ for $\ell < 10$. *Right:* Evaluating the source term (4.20) subject to $i + \gamma \geq \ell$ for the same values of χ and the same range of ℓ .

Unlike the case with all normalizable modes where two of the three resonance channels naturally vanished, both of the resonant channels contribute when the non-normalizable modes have frequencies given by (4.16). In figure 6, we evaluate both resonant contributions channels and plot their contributions for various values of χ . In particular, we examine the values $\chi \in \{\pi/6, \dots, 7\pi/6\}$.

5 Discussion

Discussion goes here.

A Derivation of Source Terms For Massive Scalars

The derivation of the general expression for the $\mathcal{O}(\epsilon^3)$ source term for massive scalars closely follows the massless case, particularly if one chooses not to write out the explicit mass dependence as was done in [5]. However, since we have chosen to write our equations in a slightly different way – and in a different gauge – than previous authors, one may find it instructive to see the differences in the derivations. Below we have included the intermediate steps involved in deriving the third-order source term S_ℓ .

Continuing the expansion of the equations of motion in powers of ϵ , we see that the back-reaction between the metric and the scalar field appears at second order in the perturbation,

$$A'_2 = -\mu\nu \left[(\dot{\phi}_1)^2 + (\phi'_1)^2 + m^2 \phi_1^2 \sec^2 x \right] + \nu' A_2 / \nu \quad (\text{A.1})$$

which can be directly integrated to give

$$A_2 = -\nu \int_0^x dy \mu \left((\dot{\phi}_1)^2 + (\phi'_1)^2 + m^2 \phi_1^2 \sec^2 x \right). \quad (\text{A.2})$$

Similarly, the first non-trivial contribution to the lapse (in the boundary time gauge) is

$$\delta_2 = \int_x^{\pi/2} dy \mu \nu \left((\dot{\phi}_1)^2 + (\phi'_1)^2 \right). \quad (\text{A.3})$$

For convenience, we have also defined the functions

$$\mu(x) = (\tan x)^{d-1} \quad \text{and} \quad \nu(x) = (d-1)/\mu'. \quad (\text{A.4})$$

To aide in evaluating integrals and inner products, it is useful to derive several identities. First, from the equation for the scalar field's time-dependent coefficients c_i ,

$$\ddot{c}_i + \omega_i^2 c_i = 0 \quad \Rightarrow \quad \partial_t (\dot{c}_i^2 + \omega_i^2 c_i^2) = \partial_t \mathbb{C}_i = 0. \quad (\text{A.5})$$

Next, from the definition of \hat{L} ,

$$\hat{L}e_j = -\frac{1}{\mu} (\mu e'_j)' + m^2 \sec^2 x e_j \quad \Rightarrow \quad (\mu e'_j)' = \mu (m^2 \sec^2 x - \omega_j^2) e_j. \quad (\text{A.6})$$

By considering the expression $(\mu e'_i e_j)'$, we see that

$$(\mu e'_i e_j)' = (m^2 \sec^2 x - \omega_i^2) \mu e_i e_j + \mu e'_i e'_j, \quad (\text{A.7})$$

which, after permuting i, j and subtracting from above, gives

$$\frac{[\mu(e'_i e_j \omega_j^2 - e_i e'_j \omega_i^2)]'}{(\omega_j^2 - \omega_i^2)} = \mu m^2 \sec^2 x e_i e_j + \mu e'_i e'_j. \quad (\text{A.8})$$

Projecting each of the terms in (2.13) individually onto the eigenbasis $\{e_\ell\}$ will involve evaluating inner products involving multiple integrals. To aide in evaluating these expressions, it is useful to derive several identities. First, from the equation for the scalar field's time-dependent coefficients c_i ,

$$\ddot{c}_i + \omega_i^2 c_i = 0 \quad \Rightarrow \quad \partial_t (\dot{c}_i^2 + \omega_i^2 c_i^2) = \partial_t \mathbb{C}_i = 0. \quad (\text{A.9})$$

Next, from the definition of \hat{L} ,

$$\hat{L}e_j = -\frac{1}{\mu} (\mu e'_j)' + m^2 \sec^2 x e_j \quad \Rightarrow \quad (\mu e'_j)' = \mu (m^2 \sec^2 x - \omega_j^2) e_j. \quad (\text{A.10})$$

By considering the expression $(\mu e'_i e_j)'$, we see that

$$(\mu e'_i e_j)' = (m^2 \sec^2 x - \omega_i^2) \mu e_i e_j + \mu e'_i e'_j, \quad (\text{A.11})$$

which, after permuting i, j and subtracting from above, gives

$$\frac{[\mu(e'_i e_j \omega_j^2 - e_i e'_j \omega_i^2)]'}{(\omega_j^2 - \omega_i^2)} = \mu m^2 \sec^2 x e_i e_j + \mu e'_i e'_j. \quad (\text{A.12})$$

Using these identities, we evaluate each of the inner products and find that

$$\begin{aligned} \langle \delta_2 \ddot{\phi}_1, e_\ell \rangle = & - \sum_{i=0}^{\infty} \sum_{\substack{j=0 \\ k \neq \ell}}^{\infty} \sum_{k=0}^{\infty} \frac{\omega_k^2 c_k}{\omega_\ell^2 - \omega_k^2} [\dot{c}_i \dot{c}_j (X_{k\ell ij} - X_{\ell k ij}) + c_i c_j (Y_{ij\ell k} - Y_{ijk\ell})] \\ & - \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \omega_\ell^2 c_\ell [\dot{c}_i \dot{c}_j P_{ij\ell} + c_i c_j B_{ij\ell}] , \end{aligned} \quad (\text{A.13})$$

$$\begin{aligned} \langle A_2 \ddot{\phi}_1, e_\ell \rangle = & 2 \sum_{i=0}^{\infty} \sum_{\substack{j=0 \\ i \neq j}}^{\infty} \sum_{k=0}^{\infty} \frac{\omega_k^2 c_k}{\omega_j^2 - \omega_i^2} X_{ijk\ell} (\dot{c}_i \dot{c}_j + \omega_j^2 c_i c_j) \\ & + \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \omega_j^2 c_j (\mathbb{C}_i P_{j\ell i} + c_i^2 X_{iij\ell}) , \end{aligned} \quad (\text{A.14})$$

$$\begin{aligned} \langle \dot{\delta}_2 \dot{\phi}_1, e_\ell \rangle = & \sum_{i=0}^{\infty} \sum_{\substack{j=0 \\ k \neq \ell}}^{\infty} \sum_{k=0}^{\infty} \frac{\dot{c}_k}{\omega_\ell^2 - \omega_k^2} [\partial_t (\dot{c}_i \dot{c}_j) (X_{k\ell ij} - X_{\ell k ij}) + \partial_t (c_i c_j) (Y_{ij\ell k} - Y_{ijk\ell})] \\ & + \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \dot{c}_\ell [\partial_t (\dot{c}_i \dot{c}_j) P_{ij\ell} + \partial_t (c_i c_j) B_{ij\ell}] , \end{aligned} \quad (\text{A.15})$$

$$\langle \dot{A}_2 \dot{\phi}_1, e_\ell \rangle = -2 \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \sum_{k=0}^{\infty} \dot{c}_k \dot{c}_j c_i X_{ijk\ell} , \quad (\text{A.16})$$

$$\begin{aligned} \langle (A'_2 - \delta'_2) \phi'_1, e_\ell \rangle = & -2 \sum_{i=0}^{\infty} \sum_{\substack{j=0 \\ i \neq j}}^{\infty} \sum_{k=0}^{\infty} \frac{c_k (\dot{c}_i \dot{c}_j + \omega_j^2 c_i c_j)}{\omega_j^2 - \omega_i^2} H_{ijk\ell} - m^2 \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \sum_{k=0}^{\infty} c_i c_j c_k V_{ijk\ell} \\ & - \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} c_j [c_i^2 H_{iij\ell} + \mathbb{C}_i M_{j\ell i}] , \end{aligned} \quad (\text{A.17})$$

$$\begin{aligned} \langle A_2 \phi_1 \sec^2 x, e_\ell \rangle = & -2 \sum_{i=0}^{\infty} \sum_{\substack{j=0 \\ i \neq j}}^{\infty} \sum_{k=0}^{\infty} \frac{c_k (\dot{c}_i \dot{c}_j + \omega_j^2 c_i c_j)}{\omega_j^2 - \omega_i^2} V_{jkil} \\ & - \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} c_j (c_i^2 V_{jii\ell} + \mathbb{C}_i Q_{j\ell i}) , \end{aligned} \quad (\text{A.18})$$

where the forms of X, Y, V, H, B, M, P, and Q are given by

$$X_{ijkl} = \int_0^{\pi/2} dx \mu^2 \nu e'_i e'_j e_k e_\ell \quad (\text{A.19})$$

$$Y_{ijkl} = \int_0^{\pi/2} dx \mu^2 \nu e'_i e'_j e_k e'_\ell \quad (\text{A.20})$$

$$V_{ijkl} = \int_0^{\pi/2} dx \mu^2 \nu e_i e_j e'_k e_\ell \sec^2 x \quad (\text{A.21})$$

$$H_{ijkl} = \int_0^{\pi/2} dx \mu^2 \nu' e'_i e'_j e'_k e_\ell \quad (\text{A.22})$$

$$B_{ij\ell} = \int_0^{\pi/2} dx \mu \nu e'_i e'_j \int_0^x dy \mu e_\ell^2 \quad (\text{A.23})$$

$$M_{ij\ell} = \int_0^{\pi/2} dx \mu \nu' e'_i e_j \int_0^x dy \mu e_\ell^2 \quad (\text{A.24})$$

$$P_{ij\ell} = \int_0^{\pi/2} dx \mu \nu e_i e_j \int_0^x dy \mu e_\ell^2 \quad (\text{A.25})$$

$$Q_{ij\ell} = \int_0^{\pi/2} dx \mu \nu e_i e_j \sec^2 x \int_0^x dy \mu e_\ell^2. \quad (\text{A.26})$$

Note that, using integration by parts to remove the derivative from ν in the definitions of H_{ijkl} and $M_{ij\ell}$, we can show that

$$H_{ijkl} = \omega_i^2 X_{kij\ell} + \omega_k^2 X_{ijkl} - Y_{ij\ell k} - Y_{\ell kji} - m^2 V_{kji\ell} - m^2 V_{ijkl}, \quad (\text{A.27})$$

$$M_{ij\ell} = \omega_i^2 P_{ij\ell} - B_{ij\ell} - m^2 Q_{ij\ell}. \quad (\text{A.28})$$

Collecting (A.13) - (A.18) gives the expression for $S_\ell = \langle S, e_\ell \rangle$:

$$\begin{aligned} S_\ell = & \sum_{\substack{i,j,k \\ k \neq \ell}}^{\infty} \frac{1}{\omega_\ell^2 - \omega_k^2} \left[F_k(\dot{c}_i \dot{c}_j) (X_{k\ell ij} - X_{\ell kij}) + F_k(c_i c_j) (Y_{ij\ell k} - Y_{ijk\ell}) \right] \\ & + 2 \sum_{\substack{i,j,k \\ i \neq j}}^{\infty} \frac{c_k D_{ij}}{\omega_j^2 - \omega_i^2} \left[2\omega_k^2 X_{ijkl} - H_{ijkl} - m^2 V_{jki\ell} \right] - \sum_{i,j,k}^{\infty} c_i \left[2\dot{c}_j \dot{c}_k X_{ijkl} + m^2 c_j c_k V_{ijkl} \right] \\ & + \sum_{i,j}^{\infty} \left[F_\ell(\dot{c}_i \dot{c}_j) P_{ij\ell} + F_\ell(c_i c_j) B_{ij\ell} + 2\omega_j^2 c_j (c_i^2 X_{iij\ell} + \mathbb{C}_i P_{j\ell i}) \right. \\ & \left. - c_j (c_i^2 (H_{iij\ell} + m^2 V_{jii\ell}) + \mathbb{C}_i (M_{j\ell i} + m^2 Q_{j\ell i})) \right], \end{aligned} \quad (\text{A.29})$$

where $F_k(z) = \dot{c}_k \dot{z} - 2\omega_k^2 c_k z$, $D_{ij} = \dot{c}_i \dot{c}_j + \omega_j^2 c_i c_j$, and $\mathbb{C}_i = \dot{c}_i^2 + \omega_i^2 c_i^2$. Additionally, we have combined some integrals into their own expressions, namely

$$Z_{ijkl}^\pm = \omega_i \omega_j (X_{k\ell ij} - X_{\ell kij}) \pm (Y_{ij\ell k} - Y_{ijk\ell}) \quad \text{and} \quad \tilde{Z}_{ij\ell}^\pm = \omega_i \omega_j P_{ij\ell} \pm B_{ij\ell}. \quad (\text{A.30})$$

Finally, using the solution for the time-dependent coefficients, $c_i(t) = a_i(t) \cos(\omega_i t + b_i(t)) \equiv a_i \cos \theta_i$, we arrive at (2.14).

B Two Non-normalizable Modes with Equal Frequencies

Let us return to the case of two, equal, non-normalizable modes with frequency $\bar{\omega}$. Within the space of resonant frequency values, there are frequencies that happen to satisfy $\bar{\omega} = \omega_\ell$ numerically and may produce extra resonances subject to restrictions on the normalizable frequency. These instances were excluded from the discussion in §4.1, and we address them here. When considering special integer values of $\bar{\omega}$ each choice of $\bar{\omega}$ below will contribute a \bar{T} -type term to the total source:

$$\bar{T}_i^{(1)} : \quad \omega_i = \omega_\ell + 2\bar{\omega} \quad \forall \bar{\omega} \in \mathbb{Z}^* \quad (\text{B.1})$$

$$\bar{T}_i^{(2)} : \quad \omega_i = \omega_\ell - 2\bar{\omega} \quad \forall \bar{\omega} \in \mathbb{Z}^* \text{ such that } \ell \geq \bar{\omega} \quad (\text{B.2})$$

$$\bar{T}_i^{(3)} : \quad \omega_i = 2\bar{\omega} - \omega_\ell \quad \forall \bar{\omega} \in \mathbb{Z}^* \text{ such that } \bar{\omega} \leq \ell + \Delta^+, \quad (\text{B.3})$$

with $\omega_i \neq \omega_\ell$ in each case. These special values contribute to the case of two, equal non-normalizable modes via

$$\begin{aligned} S_\ell = & \bar{A}_\omega^2 \bar{T}_{(\ell+\bar{\omega})}^{(1)} a_{(\ell+\bar{\omega})} \cos(\theta_{(\ell+\bar{\omega})} - 2\bar{\omega}t) + \bar{A}_\omega^2 \bar{T}_{(\ell-\bar{\omega})}^{(2)} a_{(\ell-\bar{\omega})} \cos(\theta_{(\ell-\bar{\omega})} + 2\bar{\omega}t) \\ & + \bar{A}_\omega^2 \bar{T}_{(\bar{\omega}-\ell-\Delta^+)}^{(3)} a_{(\bar{\omega}-\ell-\Delta^+)} \cos(2\bar{\omega}t - \theta_{(\bar{\omega}-\ell-\Delta^+)}) \end{aligned} \quad (\text{B.4})$$

under their respective conditions on the value of $\bar{\omega}$. The total resonant contribution for all possible $\bar{\omega}$ values is the addition of (B.4) and (4.6). Evaluating (2.14) in each case of the cases described by (B.1)-(B.3), we find that

$$\begin{aligned} \bar{T}_i^{(1)} = & \frac{1}{2} \left[H_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) - H_{\bar{\omega}i\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i - \bar{\omega}} \right) + m^2 V_{\bar{\omega}\bar{\omega}i\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) \right. \\ & \left. - m^2 V_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i - \bar{\omega}} \right) - 2\bar{\omega}^2 X_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) + 2\bar{\omega}^2 X_{\bar{\omega}i\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i - \bar{\omega}} \right) \right]_{\omega_i \neq \bar{\omega}} \\ & - \frac{1}{2} \left[Z_{i\bar{\omega}\bar{\omega}\ell}^+ \left(\frac{\bar{\omega}}{\omega_\ell + \bar{\omega}} \right) \right]_{\omega_\ell \neq \bar{\omega}} + \frac{1}{4} Z_{\bar{\omega}\bar{\omega}i\ell}^- \left(\frac{\omega_\ell + 2\bar{\omega}}{2\bar{\omega}} \right) + \frac{1}{2} \bar{\omega}^2 X_{i\bar{\omega}\bar{\omega}\ell} - \frac{m^2}{4} V_{\bar{\omega}\bar{\omega}i\ell} \\ & - \bar{\omega} \omega_i X_{\bar{\omega}\bar{\omega}i\ell} - \frac{m^2}{2} V_{i\bar{\omega}\bar{\omega}\ell}, \end{aligned} \quad (\text{B.5})$$

$$\begin{aligned} \bar{T}_i^{(2)} = & -\frac{1}{2} \left[H_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\bar{\omega}}{\omega_i + \bar{\omega}} \right) + H_{\bar{\omega}i\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i + \bar{\omega}} \right) + m^2 V_{\bar{\omega}\bar{\omega}i\ell} \left(\frac{\bar{\omega}}{\omega_i + \bar{\omega}} \right) \right. \\ & \left. + m^2 V_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i + \bar{\omega}} \right) - 2\bar{\omega}^2 X_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\bar{\omega}}{\omega_i + \bar{\omega}} \right) - 2\bar{\omega}^2 X_{\bar{\omega}i\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i + \bar{\omega}} \right) \right]_{\omega_i \neq \bar{\omega}} \\ & - \frac{1}{2} \left[Z_{i\bar{\omega}\bar{\omega}\ell}^- \left(\frac{\bar{\omega}}{\omega_\ell - \bar{\omega}} \right) \right]_{\omega_\ell \neq \bar{\omega}} - \frac{1}{4} Z_{\bar{\omega}\bar{\omega}i\ell}^- \left(\frac{\omega_\ell - 2\bar{\omega}}{\bar{\omega}} \right) + \frac{1}{2} \bar{\omega}^2 X_{i\bar{\omega}\bar{\omega}\ell} + \frac{m^2}{4} V_{\bar{\omega}\bar{\omega}i\ell} \\ & + \bar{\omega} \omega_i X_{\bar{\omega}\bar{\omega}i\ell} + \frac{m^2}{2} V_{i\bar{\omega}\bar{\omega}\ell}, \end{aligned} \quad (\text{B.6})$$

and

$$\begin{aligned}
\bar{T}_i^{(3)} = & \frac{1}{2} \left[H_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) - H_{\bar{\omega}i\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i - \bar{\omega}} \right) + m^2 V_{\bar{\omega}\bar{\omega}i\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) \right. \\
& - m^2 V_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\omega_i}{\omega_i - \bar{\omega}} \right) - 2\bar{\omega}^2 X_{i\bar{\omega}\bar{\omega}\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) + 2\omega_i^2 X_{\bar{\omega}\bar{\omega}i\ell} \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) \\
& \left. - Z_{i\bar{\omega}\bar{\omega}\ell}^+ \left(\frac{\bar{\omega}}{\omega_i - \bar{\omega}} \right) \right]_{\omega_i \neq \bar{\omega}} + \frac{1}{4} Z_{\bar{\omega}\bar{\omega}i\ell}^- \left(\frac{2\bar{\omega} - \omega_\ell}{2\bar{\omega}} \right) + \frac{1}{2} \bar{\omega}^2 X_{i\bar{\omega}\bar{\omega}\ell} - \frac{m^2}{4} V_{\bar{\omega}\bar{\omega}i\ell} \\
& - \bar{\omega} \omega_i X_{\bar{\omega}\bar{\omega}i\ell} - \frac{m^2}{2} V_{i\bar{\omega}\bar{\omega}\ell}.
\end{aligned} \tag{B.7}$$

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