

# Parameter Forecasts From CMB Lensing and Galaxy Lensing Power- and Bispectra

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**Abstract**

# 1 Introduction

Weak lensing has been used as a probe to constrain cosmological parameters since the 2000's [18, 3, 16, 27, 5]. A range of upcoming surveys, such as the survey by the Simons Observatory (SO) [1], the Legacy Survey of Space and Time (LSST) [14], and Euclid [19] aim to measure weak lensing of the CMB and of galaxies. They will do so much more accurately than previous surveys, enabling tighter constraints on parameters and progressing our understanding of physics as a whole. Typically, the main value of interest is the power spectrum (or equivalently, two point function) of the lensing potential. This can be measured in terms of the lensing convergence  $\kappa$ , lensing shear  $\gamma$ , or lensing potential  $\psi$ . In the weak lensing regime they all contain the same amount of information and can be easily converted into one another, see section B. For the CMB the lensing potential is generally measured directly [23], while with galaxies one instead measures the shear, as this can be deduced from the ellipticity distribution of the observed galaxies [12].

Even with Gaussian initial conditions, there are significant amounts of non-Gaussianity in the weak lensing signal, especially in galaxy lensing due to it being a probe of the more recent distribution of matter [6, 24]. In addition to looking at the power spectrum, a natural next step is thus to measure the lensing bispectrum. This has already been done for signals from galaxies [28], but no proper detection has been done of the CMB lensing bispectrum. Upcoming surveys such as the ones mentioned earlier may detect this signal. Therefore it is relevant to know the detectability of, and parameter constraints possible from, all 4 signals: CMB lensing power- and bispectra and galaxy lensing power- and bispectra. Notably, this includes looking at cross correlation as well.

In this article we aim to answer this question in the context of experimental parameters similar to those of next generation surveys such as mentioned earlier. We are especially interested in seeing if approximate parameter degeneracies can be broken by combining all data. Calculating the covariance matrix of unbiased optimal estimators of the cosmological parameters can be done by combining cosmological simulations (in our case we use the CAMB package [20]) with a Fisher matrix analysis. See figure 1 for an overview of the calculation done in this paper. The parameters looked at are the Hubble constant,  $H_0$ , physical baryon density,  $\Omega_b h^2$ , physical cold dark matter density,  $\Omega_c h^2$ , scalar spectral index,  $n_s$ , amplitude of primordial scalar fluctuations,  $A_s$ , sum of neutrino masses  $\sum m_\nu$ , and dark energy equation of state parameter,  $w_0$ . Though, we will find that some of the parameters are poorly constrained even when combining all 4 signals and thus omit them from the analysis in the results. As a secondary purpose, this paper also includes a derivation for the Fisher matrix formalism in the non-trivial context of bispectra with multiple tracers. The final formula can already be found in the literature (see e.g. the appendix of [17]), but the authors of this paper were not able to find a derivation and hence included it.

The structure of this paper is as follows: we first review the theory of weak lensing and weak lensing statistics in sections A and B. We then provide technical details of the calculation in section 3, such as choices for noise and numerical approximations made in the code. In section 4 we present the results, consisting of signal to noise ratios and parameter constraints for different combinations of the CMB and galaxy lensing power- and bispectra. We conclude in section 5. Appendix C contains a short review of the Fisher matrix formalism and gives a technical derivation of the formula used to calculate each Fisher matrix element in the case of bispectra measurements with multiple tracers. Finally, appendix D explains how to relate the shear to the lensing potential in spherical harmonic space.

All code written for this project is uploaded on github at [https://github.com/Jonas-Frugte/fisher\\_calc\\_weak\\_lensing](https://github.com/Jonas-Frugte/fisher_calc_weak_lensing).



Figure 1: Flowchart of steps involved in calculating the minimum error/uncertainty in cosmological parameter estimates using convergence and shear power- and bi-spectra.

## 2 Theory

### 2.1 Weak lensing spectra

Radiation from cosmological objects is distorted due to gravitational lensing. In the vast majority of cases this lensing is a relatively weak effect and is thus referred to as weak lensing. Weak lensing is quantified through the deflection field  $\mathbf{d}(\hat{\mathbf{n}})$  which equals the difference between the observed angle of a point on the sky and the true (unlensed) angle. This field is the gradient of the lensing potential  $\psi(\hat{\mathbf{n}})$  given as a weighted integral over the mass distribution along the line of sight between earth and the observed point.

In the case of CMB surveys, lensing alters the statistical properties of the temperature and polarization fields and can thus be calculated by comparing to the expected unlensed signal, see e.g. [21]. In galaxy surveys, lensing alters the ellipticities of observed galaxies. If a large enough number of galaxies are observed, this effect can be separated from the intrinsic ellipticities of the galaxies and thus the lensing potential can be estimated.

To constrain cosmological parameters we can then look at the lensing potential of the CMB,  $\psi_{\text{CMB}}$ , and of galaxy surveys,  $\psi_{\text{gal}}$ . These are directly related to the matter power and bispectra as

$$\begin{aligned} P_l^{\psi_X \psi_Y} &= \frac{9}{l^4} \Omega_m^2 H_0^4 \int_0^{\chi_*} \chi^2 d\chi a(\eta_0 - \chi)^{-2} W_X(\chi) W_Y(\chi) P^\delta(l/\chi, \eta_0 - \chi), \\ B_{l_1 l_2 l_3}^{\psi_X \psi_Y \psi_Z} &= \sqrt{\frac{(2l_1 + 1)(2l_2 + 1)(2l_3 + 1)}{4\pi}} \begin{pmatrix} l_1 & l_2 & l_3 \\ 0 & 0 & 0 \end{pmatrix} \frac{27}{l_1^2 l_2^2 l_3^2} \Omega_m^3 H_0^6 \\ &\quad \times \int \chi^2 d\chi a(\eta_0 - \chi)^{-3} W_X(\chi) W_Y(\chi) W_Z(\chi) B^\delta(\{l_i/\chi\}, \eta_0 - \chi), \end{aligned}$$

with  $X, Y, Z \in \{\text{CMB}, \text{gal}\}$ .  $\Omega_m$  is the present day matter density parameter,  $H_0$  is the present day Hubble constant,  $a(\eta)$  is the scale factor,  $\eta_0$  is the conformal time today,  $\chi$  is the comoving radial distance,  $\chi_*$  the distance to surface of last scattering,  $P^\delta(k, \eta)$  the matter power spectrum,  $B^\delta(k_1, k_2, k_3, \eta)$  the matter bispectrum,  $W_X(\chi)$  the window function, and

$$\begin{pmatrix} l_1 & l_2 & l_3 \\ m_1 & m_2 & m_3 \end{pmatrix}$$

the Wigner 3-j symbol. The only factor that differs between CMB and galaxy lensing is the window function, defined as

$$W_X(\chi) = \int_\chi^\infty d\chi' p_X(\chi') \frac{\chi' - \chi}{\chi' \chi},$$

with  $p_X$  the radius distribution of the radiation source. For the CMB we take  $p(\chi) = \delta(\chi - \chi_*)$ . Notice that we can also look at cross correlation between CMB and galaxy lensing.

For a derivation of the lensing potential and its power- and bispectra, see appendices A and B.

## 2.2 Nonlinear matter powerspectrum

Determining the nonlinear matter powerspectrum can be done using numerical codes such as CAMB. The nonlinear matter bispectrum was calculated from the powerspectrum using a fitting formula based on perturbation theory in [11]. It is given by

$$B_m(k_1, k_2, k_3, \chi) = 2F_2(k_1, k_2, z)P_m(k_1, z)P_m(k_2, z) + 2 \text{ perm.} \quad (1)$$

where  $P_m$  is the nonlinear matter power spectrum<sup>1</sup>, and the kernel  $F_2$  is modified from the tree level result with factors  $a(k, z)$ ,  $b(k, z)$ , and  $c(k, z)$ :

$$F_2(k_1, k_2, z) = \frac{5}{7}a(k_1, z)a(k_2, z) + \frac{k_1^2 + k_2^2}{2k_1k_2}b(k_1, z)b(k_2, z)\cos\theta + \frac{2}{7}c(k_1, z)c(k_2, z)\cos^2\theta. \quad (2)$$

They are defined as:

$$a(k, z) = \frac{1 + \sigma_8^{a_6}(z)\sqrt{0.7}Q(n_{\text{eff}})(q^{a_1})^{n_{\text{eff}}+a_2}}{1 + (q^{a_1})^{n_{\text{eff}}+a_2}}, \quad (3)$$

$$b(k, z) = \frac{1 + 0.2a_3(n_{\text{eff}} + 3)(q^{a_7})^{n_{\text{eff}}+3+a_8}}{1 + (q^{a_5})^{n_{\text{eff}}+3.5+a_8}}, \quad (4)$$

$$c(k, z) = \frac{1 + \left[ \frac{4.5a_4}{1.5 + (n_{\text{eff}}+3)^4} \right] (q^{a_5})^{n_{\text{eff}}+3+a_9}}{1 + (q^{a_5})^{n_{\text{eff}}+3.5+a_9}}. \quad (5)$$

Here,  $Q(n_{\text{eff}})$  is given by:

$$Q(x) = \frac{4 - 2x}{1 + 2^{x+1}}. \quad (6)$$

The effective spectral index of the linear power spectrum is defined as:

$$n_{\text{eff}} \equiv \frac{d \ln P_m^{\text{lin}}(k)}{d \ln k}. \quad (7)$$

Additionally,  $q$  is given by:

$$q = \frac{k}{k_{\text{NL}}}, \quad (8)$$

where  $k_{\text{NL}}$  is the scale at which nonlinearities become significant, satisfying:

$$4\pi k_{\text{NL}}^3 P_m^{\text{lin}}(k_{\text{NL}}, 0) = 1. \quad (9)$$

The coefficients  $a_i$  are:

$$\begin{aligned} a_1 &= 0.484, & a_2 &= 3.740, & a_3 &= -0.849, & a_4 &= 0.392, \\ a_5 &= 1.013, & a_6 &= -0.575, & a_7 &= 0.128, & a_8 &= -0.722, & a_9 &= -0.926. \end{aligned}$$

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<sup>1</sup>Compare to the tree level bispectrum where we instead use the linear powerspectrum.

### 2.3 Fisher matrix analysis

SNR and parameter constraints are obtained through a Fisher matrix analysis as is standard. A full derivation of below statements can be found in appendix C. The auto and cross power spectrum Fisher matrix is given as

$$F_{\alpha\beta} = \sum_l \sum_{XY} \sum_{X'Y'} (2l+1) \partial_\alpha C_l^{XY} (C^{-1})_l^{XX'} (C^{-1})_l^{YY'} \partial_\beta C_l^{X'Y'}$$

where

$$C_l := \begin{pmatrix} C_l^{\psi_\kappa\psi_\kappa} & C_l^{\psi_\kappa\psi_\gamma} \\ C_l^{\psi_\kappa\psi_\gamma} & C_l^{\psi_\gamma\psi_\gamma} \end{pmatrix}.$$

For the bispectra we instead get

$$F_{\alpha\beta} = \sum_{l_1 \leq l_2 \leq l_3} \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \sum_{XYZ} \sum_{X'Y'Z'} \partial_\alpha B_{l_1 l_2 l_3}^{XYZ} (C^{-1})_{l_1}^{XX'} (C^{-1})_{l_2}^{YY'} (C^{-1})_{l_3}^{ZZ'} \partial_\beta B_{l_1 l_2 l_3}^{X'Y'Z'},$$

with  $\mathcal{P}_{l_1 l_2 l_3}$  defined as the number of distinct permutations that can be made with  $l_1 l_2 l_3$

When only considering autospectra the formulas simplify to

$$F_{\alpha\beta} = \sum_l \sum_{XY} \sum_{X'Y'} \frac{2l+1}{2} \frac{\partial_\alpha C_l^{XX} \partial_\beta C_l^{XX}}{(C_l^{XX})^2},$$

and, for bispectra,

$$F_{\alpha\beta} = \sum_{l_1 \leq l_2 \leq l_3} \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \frac{\partial_\alpha B_{l_1 l_2 l_3}^{XXX} \partial_\beta B_{l_1 l_2 l_3}^{XXX}}{C_{l_1}^{XX} C_{l_2}^{XX} C_{l_3}^{XX}}.$$

## 3 Calculation Details

We chose our fiducial cosmology in line with the results from the Planck mission [2].

$H_0 = 67.4 \text{ km/s/Mpc},$	(Hubble constant)
$\Omega_b h^2 = 0.0223,$	(Physical baryon density parameter)
$\Omega_c h^2 = 0.119,$	(Physical cold dark matter density parameter)
$n_s = 0.965,$	(Scalar spectral index)
$A_s = 2.13 \times 10^{-9},$	(Amplitude of primordial scalar fluctuations)
$\sum m_\nu = 0.06 \text{ eV},$	(Sum of neutrino masses)
$w_0 = -1.$	(Dark energy equation of state parameter)

This paper considers noise levels for stage 3 and stage 4 weak lensing surveys. Values can be found in table 1. A comparison of noise power spectrum to lensing power spectrum can be found in figure 2.

source	survey stage	noise vals	comparable experiments
CMB	stage 3	$\sigma = 1', \Delta_P = 6' \mu\text{K}$	Advanced ACTPol, SPT-3G
	stage 4	$\sigma = 3', \Delta_P = 1' \mu\text{K}$	CMB-S4, LiteBIRD
galaxies	stage 3	$\sigma_{\text{rms}} = 0.3, n_g = 5 \text{ arcmin}^{-2}$	DES, KiDS
	stage 4	$\sigma_{\text{rms}} = 0.3, n_g = 30 \text{ arcmin}^{-2}$	LSST, Euclid

Table 1: Noise levels considered for weak lensing of galaxies and the CMB.  $\sigma$  (beam width) and  $\Delta_P$  (polarization white noise) describe CMB survey specifications, while  $\sigma_{\text{rms}}$  (intrinsic galaxy ellipticity) and  $n_g$  (observed galaxy density) refer to galaxy shear surveys.

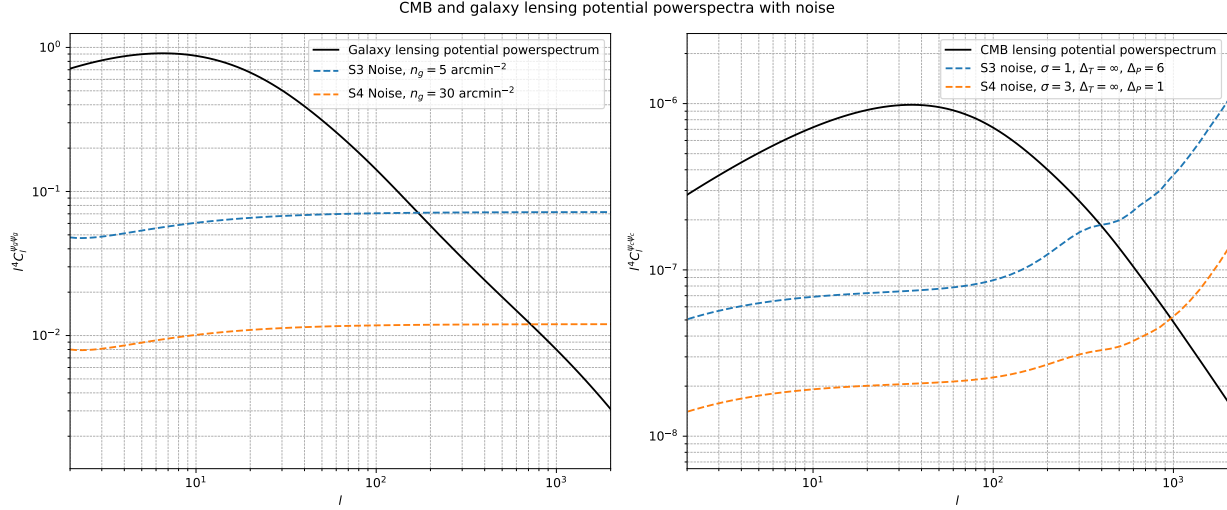


Figure 2: CMB (right) and galaxy (left) lensing potential power spectra compared to associated experimental noise. Current (stage 3) noise values are displayed as well as near future (stage 4) noise values. The CMB lensing experiment does not measure temperature anisotropies, only polarization. CMB noise values are chosen in accordance with [22]. Shear noise values are chosen to be similar to e.g. Euclid measurements for stage 4 and e.g. KiDS for stage 3.

CMB noise is estimated through the quadratic estimator developed in [21]. We do not take temperature measurements into account, as these are subleading for stage 4 experiments (REF). This is similarly done in [22]. The parameters characterizing the noise levels are thus beam width,  $\sigma$  and polarization white noise,  $\Delta_P$ . Galaxy lensing is determined by measuring lensing shear (see appendices A and D). The noise in this measurement is dominated by scale independent shot noise and has associated noise power spectrum  $N_l^{\text{shear}} = \sigma_{\text{rms}}^2 / n_g$ .  $n_g$  is the amount of galaxies observed per unit angle. The noise power spectrum for the lensing potential then equals

$$N_l^{\text{lens. potential}} = \frac{(l-1)l(l+1)(l+2)}{4} N_l^{\text{shear}}.$$

For parameter constraints we use stage 4 noise levels.

The redshift distribution of the observed galaxies is commonly parameterized as [4]

$$n(z) \propto z^a \exp \left[ - \left( \frac{z}{z_0} \right)^b \right].$$

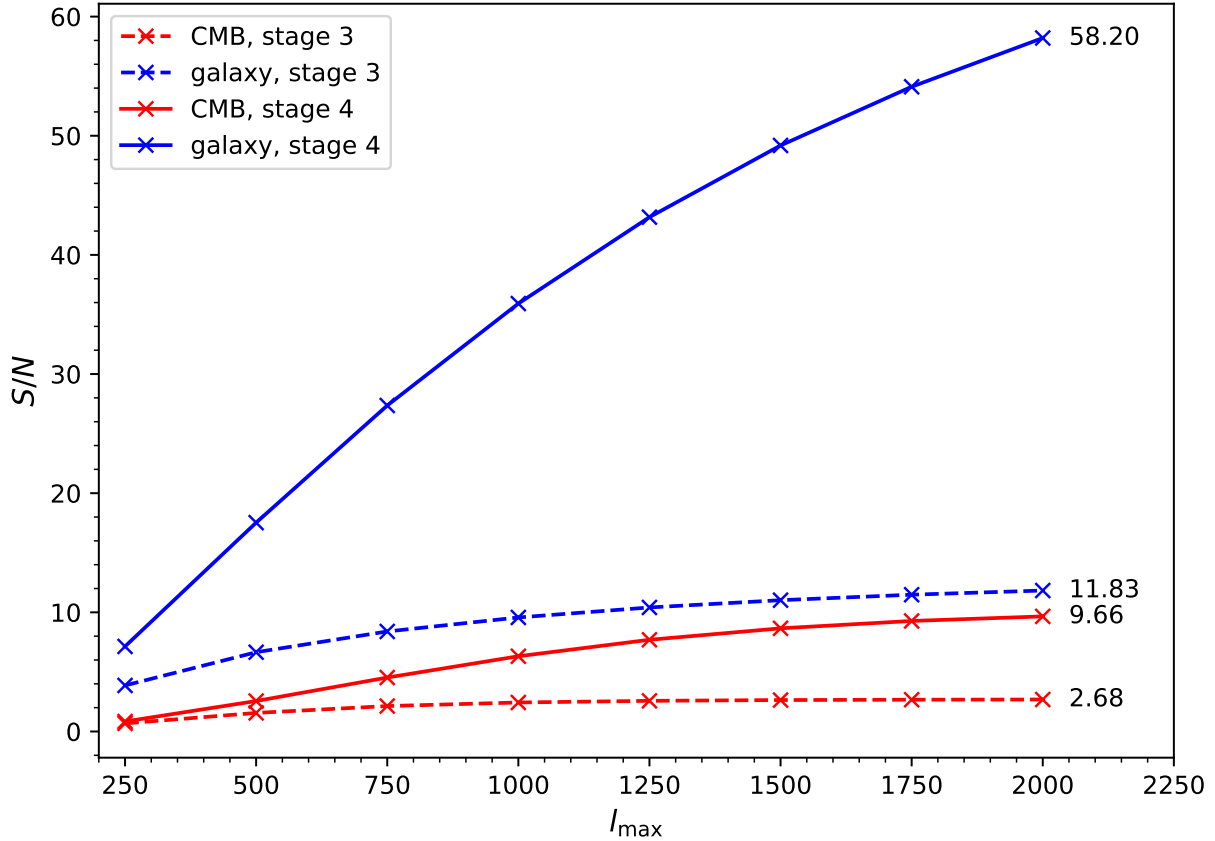


Figure 3: Signal to noise ratios for galaxy lensing (dashed) and CMB lensing (non-dashes) power- (in red) and bispectra (in blue). The  $x$ -axis labels the maximum multipole moment measured up to, starting from  $l_{\min} = 2$ .

We choose the parameter combination  $a = 2$ ,  $b = 3/2$  and  $z_0 = 0.64$  which is similar to the expected distributions of Euclid, (which will probe primarily in the 0.2 - 2.6 redshift range [9]) and the LSST mission (which has  $a \approx 2$ ,  $b \approx 1$ , and  $z_0 \approx 0.3$  from predictions for the obtained data [15]).

## 4 Results

### 4.1 SNR

Lensing power spectra are well detectable in nearly all cases. The signal to noise ratios of the CMB and galaxy lensing bispectra v.s. the maximum multipole moment measured can be found in figure 3. Shear bispectra can be measured well by both stage 3 and stage 4 experiments, while CMB lensing is only semi detectable by a stage 3 experiment but well detectable by stage 4 experiments.

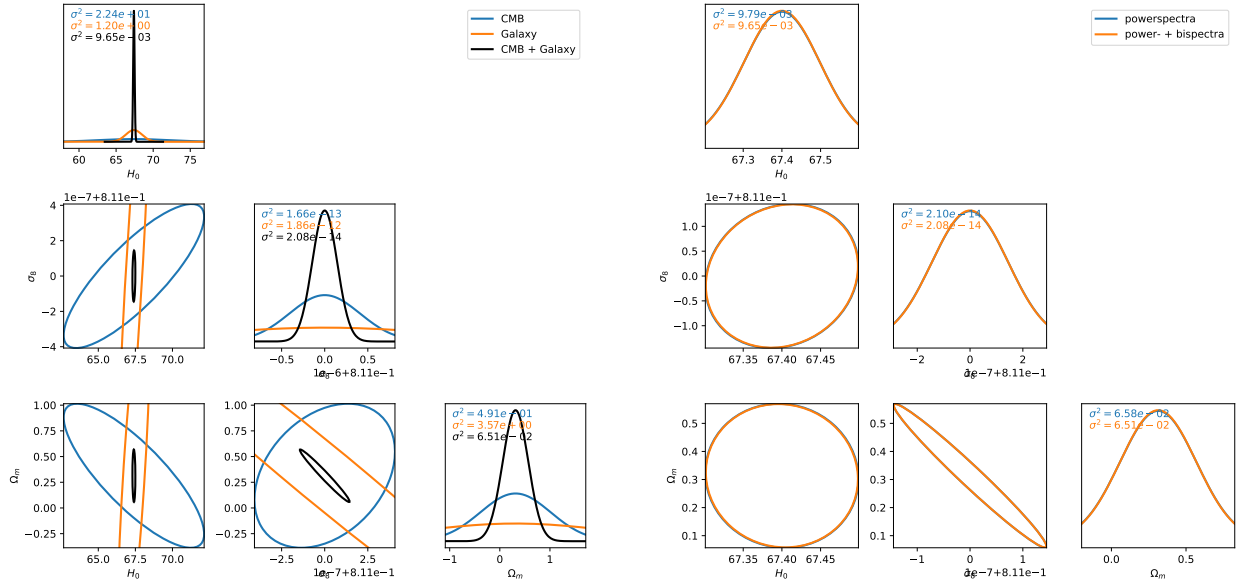


## 4.2 Parameter Constraints

For the parameter constraints, even when combining all signals, most of the parameters are poorly constrained (with minimum uncertainties much bigger than their fiducial values). The combination that is well constrained is  $H_0$  and  $\sigma_8$ . Alternatively one can also look at  $H_0$  and  $S_8$ , where  $S_8$  is defined as

$$S_8 = \sigma_8 \sqrt{\frac{\Omega_c + \Omega_b}{0.3}}$$

and is commonly found to be a particularly well constrained parameter in weak lensing surveys. The results are essentially the same so we will stick with  $\sigma_8$ . The constraints and covariances are shown in figures 4 (a) and (b). In particular, you can see in the off diagonal plot that combining the CMB and galaxy lensing removes approximate degeneracies in  $H_0$ ,  $\sigma_8$  space to create a much smaller  $1\sigma$  confidence ellipse. When comparing powerspectra to bispectra constraints in figure 4b we see a more mild improvement in constraints when combining the different sources of information.



(a) Constraints for  $H_0$  and  $\sigma_8$  for CMB lensing, galaxy lensing, and CMB + galaxy lensing combined. We use the information from both power and bispectra together with  $l_{\max} = 2000$  in all three cases. The confidence ellipses are for  $1\sigma$  ( $\sim 70\%$ ) certainty.

(b) Same as Fig. 4a, except comparing constraints from only using power spectra, only bispectra, or both.  $S_8$  instead of  $\sigma_8$ . In all three cases we use CMB and galaxy lensing.

Figure 4: Side-by-side comparison of parameter constraints.

## 5 Discussion and conclusion

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## A Weak Lensing

### A.1 Approximations Made / Regime

### A.2 Perturbed Photon Paths

Work in **conformal newtonian gauge**. In this section we’ll work in units where  $c = 1$ . Denoting conformal time and conformal radial distance by  $\eta$  and  $\chi$ , respectively, the perturbed line element

is given by

$$ds^2 = a^2(\eta)((1 + 2\Psi_N)d\eta^2 - (1 + 2\Phi_N)\gamma_{ij}dx^i dx^j) \quad (10)$$

where  $\gamma_{ij}$  is the unperturbed line element

$$\gamma_{ij} = dx^i dx^j = d\chi^2 + f_K^2(\chi)(d\theta^2 + \sin^2\theta d\phi^2). \quad (11)$$

From here of on work in first order in the scalar potentials  $\Phi_N$  and  $\Psi_N$ . For lensing we consider null-geodesics so  $ds^2 = 0$  and we can rewrite the perturbed line element as

$$d\hat{s}^2 = (1 + 4\Psi)d\eta^2 - \gamma_{ij}dx^i dx^j \quad (12)$$

with  $\Psi$  the **Weyl Potential** given by  $\Psi := (\Psi_N - \Phi_N)/2$ . A null geodesic  $x^\mu(\hat{\lambda})$  in terms of its affine parameter  $(\hat{\lambda})$  satisfies the geodesic equation

$$\frac{d^2 x^\mu}{d\hat{\lambda}^2} + \Gamma_{\nu\rho}^\mu \frac{dx^\nu}{d\hat{\lambda}} \frac{dx^\rho}{d\hat{\lambda}} = 0, \quad (13)$$

with  $\hat{g}_{\mu\nu}(dx^\mu/d\hat{\lambda})(dx^\nu/d\hat{\lambda}) = 0$ .

In the unperturbed there are incoming radial solutions of the form  $\chi = \eta_0 - \eta$ . We can express  $\hat{\lambda}$  in terms of  $\eta$  using The 0-component of the geodesic equation,

$$\frac{d^2 \eta}{d\hat{\lambda}^2} + 2\left(\frac{d\eta}{d\hat{\lambda}}\right)^2 \frac{d\Psi}{d\eta} + 2\frac{d\eta}{d\hat{\lambda}} \frac{dx^i}{d\hat{\lambda}} \frac{\partial\Psi}{\partial x^i} = 0, \quad (14)$$

where the derivative  $d\Psi/d\eta = \partial_\eta \Psi + (dx^i/d\eta)\partial_i \Psi$  is along the perturbed ray. This allows us to rewrite the other (spatial) geodesic equations as

$$\frac{d^2 x^i}{d\eta^2} - 2\frac{dx^i}{d\eta} \left( \frac{d\Psi}{d\eta} + \frac{dx^j}{d\eta} \frac{\partial\Psi}{\partial x^j} \right) + 2\gamma^{ij} \frac{\partial\Psi}{\partial x^j} + \bar{\Gamma}_{jk}^i \frac{dx^j}{d\eta} \frac{dx^k}{d\eta} = 0, \quad (15)$$

where  $\bar{\Gamma}_{jk}^i$  are the connection coefficients of the unperturbed three geometry  $\gamma_{ij}$ . Without loss of generality we consider an observer located at the origin of the spatial coordinates, meaning we are interested in rays that end at  $x^i = 0$ . In that case  $d\chi/d\eta = -1 + O(\Psi)$ ,  $d\theta/d\eta = O(\Psi)$  and  $d\phi/d\eta = O(\Psi)$  (with  $\theta$  and  $\phi$  angular spatial coordinates). For the case of such rays we can evaluate the connection coefficients and rewrite the spatial geodesic equations as

$$\frac{\partial^2 \chi}{d\eta^2} + 2\frac{d\Psi}{d\eta} = 0, \quad (16)$$

$$\frac{d^2 \theta}{d\eta^2} - 2\frac{d \ln f_K(\chi)}{d\chi} \frac{d\theta}{d\eta} + \frac{2}{f_K^2(\chi)} \frac{\partial\Psi}{\partial \theta} = 0, \quad (17)$$

$$\frac{d^2 \phi}{d\eta^2} - 2\frac{d \ln f_K(\chi)}{d\chi} \frac{d\phi}{d\eta} + \frac{2}{f_K^2(\chi)} \frac{1}{\sin^2(\theta)} \frac{\partial\Psi}{\partial \phi} = 0. \quad (18)$$

Integrating twice and using the null condition we can rewrite the the first equation as

$$\chi = \eta_0 - \eta - 2 \int_{\eta_0}^{\eta} \Psi d\eta', \quad (19)$$

where the integral is along the ray. Since we are working at first order in  $\Psi$  we can evaluate the integral along the unperturbed path, this is known as the **Born approximation**. *Next follows a short discussion of why perturbations in observation time can be ignored, include?* Doing the same for the other two equations gives

$$\theta(\eta_0 - \chi_*) = \theta_0 - \int_0^{\chi_*} d\chi \frac{f_K(\chi_* - \chi)}{f_K(\chi_*)f_K(\chi)} 2 \frac{\partial}{\partial \theta} \Psi(\chi \hat{n}, \eta_0 - \chi), \quad (20)$$

$$\phi(\eta_0 - \chi_*) = \theta_0 - \int_0^{\chi_*} d\chi \frac{f_K(\chi_* - \chi)}{f_K(\chi_*)f_K(\chi)} \frac{2}{\sin^2 \theta} \frac{\partial}{\partial \phi} \Psi(\chi \hat{n}, \eta_0 - \chi), \quad (21)$$

where  $\theta_0$  and  $\phi_0$  label the line of sight  $\hat{n}$ . The displacement vector  $\alpha$  indicates the infinitesimal displacement of the lensed rays from their unperturbed observed angles and is related to  $\theta$  and  $\phi$  through

$$\alpha_\theta = \theta - \theta_0, \quad \alpha_\phi = \sin^2 \theta (\phi - \phi_0).$$

This allows us to summarize the above results as

$$\alpha = -2 \int_0^{\chi_*} d\chi \frac{f_K(\chi_* - \chi)}{f_K(\chi_*)f_K(\chi)} \nabla_{\hat{n}} \Psi(\chi \hat{n}; \eta_0 - \chi). \quad (22)$$

When working in the weak field limit (i.e. small spacetime curvature) the 00 component of Einsteins equations give us a poisson equation [10],

$$-\nabla^2 \Psi(\mathbf{x}, \eta) = 4\pi G a(\eta)^2 \delta \rho(\mathbf{x}, \eta) \quad (23)$$

where  $K$  is 0, 1, or  $-1$  depending on the curvature of the universe,  $\Delta = \partial_x^2 + \partial_y^2 + \partial_z^2$  is the 3 dimensional laplace operator,  $\delta \rho$  is the comoving total energy perturbation (evaluated in the rest-frame of the total energy).

### A.3 Convergence, shear

Quantities we are interested in measuring depend on how the lensing vector changes with the observation angle. **Magnification matrix** is defined by

$$A_{ij} := \delta_{ij} + \frac{\partial}{\partial \theta_i} \alpha_j(\hat{n}). \quad (24)$$

This matrix can then be decomposed in the following form, which immediately gives us definitions for the **convergence**,  $\kappa$ , **shear**,  $\gamma_1$  and  $\gamma_2$ , and **rotation**,  $\omega$ :

$$A_{ij}(\hat{n}) = \begin{pmatrix} 1 - \kappa - \gamma_1 & -\gamma_2 + \omega \\ -\gamma_2 - \omega & 1 - \kappa + \gamma_1 \end{pmatrix}. \quad (25)$$

At 1st order  $A$  is a symmetric matrix by definition and so  $\omega$  has to be 0, we will ignore it from here on out. Intuitively,  $A$  tells you how a small patch in the sky transforms due to weak lensing, i.e. if you have surface brightness  $I$  then

$$I_{\text{emitted}}(\hat{n} + \delta) = I_{\text{observed}}((\hat{n} + \delta) + \alpha(\hat{n} + \delta)) \approx I_{\text{observed}}((\hat{n} + \alpha + A_{ij} \delta_j)). \quad (26)$$

given that  $|A_{ij}| = (1 - \kappa)^2 + \omega^2 - |\gamma|^2 = 1 - 2\kappa + O(\kappa^2, \gamma^2, \omega^2)$ , we can interpret  $\kappa$  as telling us about the overall magnification of the source. The  $\gamma_i$  represent the area preserving distortion, i.e. stretching and squeezing in a specific direction.

## A.4 lensing potential

Taking the derivative out of the integral in equation REF leads us to defining the **lensing potential**,  $\psi$ ,

$$\alpha = \nabla_{\hat{\mathbf{n}}} \psi, \quad \psi(\hat{\mathbf{n}}) = -2 \int_0^{\chi_*} d\chi \frac{f_K(\chi_* - \chi)}{f_K(\chi_*) f_K(\chi)} \Psi(\chi \hat{\mathbf{n}}; \eta_0 - \chi). \quad (27)$$

We can relate  $\kappa$  and  $\gamma$  directly to the lensing potential as

$$\kappa = \frac{1}{2} \nabla^2 \psi, \quad \gamma_1 = \frac{1}{2} (\partial_\theta^2 - \partial_\phi^2) \psi, \quad \gamma_2 = \partial_\theta \partial_\phi \psi. \quad (28)$$

## A.5 Regarding window functions

The window function for light emitted from a source at fixed radius  $\chi_*$  is  $(\chi - \chi_*)/(\chi_* \chi)$ , if the source is instead distributed over the radius as  $p(\chi)$ , with  $p(\chi)$  normalized to integrate to 1, we generalize the window function as

$$W(\chi) = \int_\chi^\infty d\chi' p(\chi') \frac{\chi' - \chi}{\chi' \chi}.$$

The integration limit is up to surface of last scattering because further than that no source can contribute to the observed image. In the case of an image of the CMB we can take  $p(\chi') = \delta(\chi' - \chi_*)$ , in which case the window function reduces back to  $(\chi_* - \chi)/(\chi_* \chi)$ . In the case of different sources, we will differentiate between their respective window functions with a subscript. In this case the factor that changes is the distribution function. The sources considered in this paper are only the CMB and selections of galaxies from galaxy surveys.

# B Weak Lensing Statistics

## B.1 Lensing Potential Powerspectrum

Consider the lensing potential,  $\psi$ , then decomposing into spherical harmonics gives

$$\psi(\hat{\mathbf{n}}) = \sum_{lm} \psi_{lm} Y_{lm}(\hat{\mathbf{n}}). \quad (29)$$

On the other hand, consider the decomposition of  $\Psi$  in fourier modes with the Fourier convention  $\Psi(\mathbf{x}, \eta) = \int \frac{d^3 \mathbf{k}}{(2\pi)^3} \Psi(\mathbf{k}, \eta) e^{i\mathbf{k} \cdot \mathbf{x}}$ ,

$$\psi(\hat{\mathbf{n}}) = -2 \int_0^{\chi_*} d\chi W(\chi) \int \frac{d^3 \mathbf{k}}{(2\pi)^3} \Psi(\mathbf{k}, \eta_0 - \chi) e^{i\mathbf{k} \cdot \hat{\mathbf{n}} \chi}. \quad (30)$$

We can than relate the multipole modes of  $\psi$  to the fourier modes of  $\Psi$  through

$$\psi_{lm} = \langle Y_l^m | \psi \rangle = \int d^2 \hat{\mathbf{n}} Y_l^m(\hat{\mathbf{n}})^* \psi(\hat{\mathbf{n}}) \quad (31)$$

$$= -2 \int d^2 \hat{\mathbf{n}} Y_l^m(\hat{\mathbf{n}})^* \int_0^{\chi_*} d\chi W(\chi) \int \frac{d^3 \mathbf{k}}{(2\pi)^3} \Psi(\mathbf{k}, \eta_0 - \chi) e^{i\mathbf{k} \cdot \hat{\mathbf{n}} \chi} \quad (32)$$

Now define the power spectrum as

$$\langle \Psi(\mathbf{k}, \eta) \Psi(\mathbf{k}', \eta') \rangle = \frac{2\pi^2}{k^3} P_\Psi(k, \eta, \eta') \delta(\mathbf{k} + \mathbf{k}'), \quad (33)$$

with  $\eta$  denoting the conformal time. This gives

$$\langle \psi(\hat{\mathbf{n}}) \psi(\hat{\mathbf{n}}') \rangle = 4 \int_0^{\chi^*} d\chi \int_0^{\chi^*} d\chi' W(\chi) W(\chi') \int \frac{d^3\mathbf{k}}{(2\pi)^6} \frac{2\pi^2}{k^3} P_\psi(k, \eta_0 - \chi, \eta_0 - \chi') e^{i\mathbf{k} \cdot \hat{\mathbf{n}}\chi} e^{-i\mathbf{k} \cdot \hat{\mathbf{n}}'\chi'}, \quad (34)$$

where we used that  $\eta = \eta_0 - \chi$  along the unperturbed photon path (Born approximation), with  $\eta_0$  the time at which the light ray hits earth. We can use the result

$$e^{i\mathbf{k} \cdot \hat{\mathbf{n}}\chi} = 4\pi \sum_{lm} i^l j_l(k\chi) Y_l^m(\hat{\mathbf{n}})^* Y_l^m(\hat{\mathbf{k}}) = 4\pi \sum_{lm} i^l j_l(k\chi) Y_l^m(\hat{\mathbf{n}}) Y_l^m(\hat{\mathbf{k}})^*, \quad (35)$$

where  $j_l$  is the spherical Bessel function, to rewrite the above equation. Using both versions of the identity above, we immediately get a factor  $Y_l^m(\hat{\mathbf{k}}) Y_{l'}^{m'}(\hat{\mathbf{k}})^*$  in our integral. We can factor the differential element of  $d^3\mathbf{k}$  into a radial and angular part as  $k^2 dk d^2\Omega_k$ , with  $\Omega_k$  the solid angle, to apply the orthonormality condition of the spherical harmonics. Note that we take the same sequence of steps a number of times in other parts of the derivations of the lensing spectra. We thus obtain

$$\langle \psi(\hat{\mathbf{n}}) \psi(\hat{\mathbf{n}}') \rangle = 4(4\pi)^2 \sum_{ll'mm'} i^{l-l'} \int_0^{\chi^*} d\chi \int_0^{\chi^*} d\chi' W(\chi) W(\chi') \quad (36)$$

$$\times \int \frac{k^2 dk}{(2\pi)^6} \frac{2\pi^2}{k^3} j_l(k\chi) j_{l'}(k\chi') P_\psi(k, \eta_0 - \chi, \eta_0 - \chi') Y_{lm}(\hat{\mathbf{n}}) Y_{l'm'}(\hat{\mathbf{n}}')^* \delta_{ll'} \delta_{mm'}. \quad (37)$$

The angular power spectrum is defined similarly to the power spectrum, i.e.

$$\langle \psi_{lm} \psi_{l'm'}^* \rangle = \delta_{ll'} \delta_{mm'} C_l^\psi. \quad (38)$$

Note that the correlation is independent of  $m$  and  $m'$ . We can thus read from equation 37 that

$$C_l^\psi = 4(4\pi)^2 \int_0^{\chi^*} d\chi \int_0^{\chi^*} d\chi' W(\chi) W(\chi') \int \frac{k^2 dk}{(2\pi)^6} \frac{2\pi^2}{k^3} j_l(k\chi) j_l(k\chi') P_\psi(k, \eta_0 - \chi, \eta_0 - \chi'), \quad (39)$$

which can be simplified to

$$C_l^\psi = \frac{2}{\pi^2} \int_0^{\chi^*} d\chi \int_0^{\chi^*} d\chi' W(\chi) W(\chi') \int k^2 dk j_l(k\chi) j_l(k\chi') \frac{P_\psi(k, \eta_0 - \chi, \eta_0 - \chi')}{k^3}. \quad (40)$$

To further evaluate the integral we will make the standard **Limber approximation**. The Bessel functions peak sharply at  $l = x^2$ , with the peak being increasingly sharp for higher  $l$ . Similarly to  $\delta(x - x_0)f(x) = \delta(x - x_0)f(x_0)$ , we thus take  $j_l(k\chi)f(k) \approx j_l(k\chi)f(l/\chi)$ . The Bessel functions satisfy an orthogonality condition,

$$\int k^2 dk j_l(k\chi) j_l(k\chi') = \frac{\pi}{2\chi^2} \delta(\chi - \chi'). \quad (41)$$

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<sup>2</sup>Some sources use  $x \approx l + 1/2$  instead, which is slightly more accurate for larger scales (low  $l$ ) and slightly less accurate for smaller scales.

In combination with the Limber approximation we thus find

$$\int k^2 dk j_l(k\chi) j_l(k\chi') f(k) \approx \frac{\pi}{2\chi^2} \delta(\chi - \chi') f(l/\chi). \quad (42)$$

It allows us to write the Limber-approximate angular spectrum as

$$C^\psi(l) = \frac{2}{\pi^2} \int_0^{\chi^*} d\chi \int_0^{\chi^*} d\chi' W(\chi) W(\chi') \frac{\pi}{2\chi^2} \delta(\chi - \chi') \frac{\chi^3}{l^3} P_\psi(l/\chi, \eta_0 - \chi, \eta_0 - \chi') \quad (43)$$

$$= \frac{1}{l^3 \pi} \int_0^{\chi^*} \chi d\chi W(\chi)^2 P_\psi(l/\chi, \eta_0 - \chi, \eta_0 - \chi). \quad (44)$$

## B.2 Lensing potential bispectrum

The derivation of the bispectrum proceeds similarly to that of the power spectrum. We aim to compute the bispectrum of the lensing potential fields of 3 (possibly distinct sources),  $\psi_1, \psi_2, \psi_3$ .

$$\begin{aligned} \langle (\psi_1)_{l_1 m_1} (\psi_2)_{l_2 m_2} (\psi_3)_{l_3 m_3} \rangle &= \prod_i \left( -2 \int d^2 \hat{\mathbf{n}}_i (Y_{l_i}^{m_i}(\hat{\mathbf{n}}_i))^* \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{d^3 \mathbf{k}_i}{(2\pi)^3} e^{i \mathbf{k}_i \cdot \hat{\mathbf{n}}_i \chi_i} \right) \\ &\quad \times \langle \prod_i \Psi(\mathbf{k}_i, \eta_0 - \chi_i) \rangle. \end{aligned}$$

Defining the bispectrum of the gravitational potential as

$$\langle \prod_{i=1,2,3} \Psi(\mathbf{k}_i, \eta_0 - \chi_i) \rangle = (2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\}),$$

we rewrite the lensing potential bispectrum as

$$\begin{aligned} \langle (\psi_1)_{l_1 m_1} (\psi_2)_{l_2 m_2} (\psi_3)_{l_3 m_3} \rangle &= \prod_i \left( -2 \int d^2 \hat{\mathbf{n}}_i (Y_{l_i}^{m_i}(\hat{\mathbf{n}}_i))^* \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{d^3 \mathbf{k}_i}{(2\pi)^3} e^{i \mathbf{k}_i \cdot \hat{\mathbf{n}}_i \chi_i} \right) \\ &\quad \times (2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\}). \end{aligned}$$

Now using equation 35 to rewrite the complex exponential:

$$\begin{aligned} &\langle (\psi_1)_{l_1 m_1} (\psi_2)_{l_2 m_2} (\psi_3)_{l_3 m_3} \rangle \\ &= \prod_i \left( -2 \int d^2 \hat{\mathbf{n}}_i (Y_{l_i}^{m_i}(\hat{\mathbf{n}}_i))^* \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{d^3 \mathbf{k}_i}{(2\pi)^3} 4\pi \sum_{lm} i^l j_l(k_i \chi_i) Y_l^m(\hat{\mathbf{n}}_i) Y_l^m(\hat{\mathbf{k}}_i)^* \right) \\ &\quad \times (2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\}) \\ &= \left[ \prod_i \left( -2 \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{d^3 \mathbf{k}_i}{(2\pi)^3} 4\pi i^{l_i} j_{l_i}(k_i \chi_i) Y_{l_i}^{m_i}(\hat{\mathbf{k}}_i)^* \right) \right] (2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\}). \end{aligned}$$

We can write the 3D dirac delta function in terms of spherical harmonics as

$$\delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) = 8 \int d^3 \mathbf{x} \prod_{i=1,2,3} \left( \sum_{l_j m_j} i^{l_j} j_{l_j}(k_i x) Y_{l_j}^{m_j}(\hat{\mathbf{k}}_i) Y_{l_j}^{m_j}(\hat{\mathbf{x}})^* \right).$$



This results in

$$\begin{aligned}
\langle (\psi_1)_{l_1 m_1} (\psi_2)_{l_2 m_2} (\psi_3)_{l_3 m_3} \rangle &= \prod_i \left( -2 \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{d^3 \mathbf{k}_i}{(2\pi)^3} 4\pi i^{l_i} j_{l_i}(k_i \chi_i) Y_{l_i}^{m_i}(\hat{\mathbf{k}}_i)^* \right) \\
&\quad \times (2\pi)^3 8 \int d^3 \mathbf{x} \prod_i \left( \sum_{lm} i^l j_l(k_i x) Y_l^m(\hat{\mathbf{k}}_i) Y_l^m(\hat{\mathbf{x}})^* \right) B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\}) \\
&= (2\pi)^3 8 \int d^3 \mathbf{x} \prod_i \left( -2 \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{k_i^2 dk_i}{(2\pi)^3} 4\pi (-1)^{l_i} j_{l_i}(k_i \chi_i) j_{l_i}(k_i x) Y_{l_i}^{m_i}(\hat{\mathbf{x}})^* \right) B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\})
\end{aligned}$$

The angular part of the  $\mathbf{x}$  integral can be evaluated using the identity

$$\begin{aligned}
&\int d\Omega_{\hat{n}} Y_{l_1 m_1}^*(\hat{x}) Y_{l_2 m_2}^*(\hat{n}) Y_{l_3 m_3}^*(\hat{n}) = (-1)^{m_1+m_2+m_3} \int d\Omega_{\hat{n}} Y_{l_1-m_1}(\hat{n}) Y_{l_2-m_2}(\hat{n}) Y_{l_3-m_3}(\hat{n}) \\
&= (-1)^{m_1+m_2+m_3} \sqrt{\frac{(2l_1+1)(2l_2+1)(2l_3+1)}{4\pi}} \begin{pmatrix} l_1 & l_2 & l_3 \\ 0 & 0 & 0 \end{pmatrix} \begin{pmatrix} l_1 & l_2 & l_3 \\ -m_1 & -m_2 & -m_3 \end{pmatrix} \equiv A_1^{\mathbf{m}},
\end{aligned}$$

giving

$$\begin{aligned}
\langle (\psi_1)_{l_1 m_1} (\psi_2)_{l_2 m_2} (\psi_3)_{l_3 m_3} \rangle &= (2\pi)^3 8 A_1^{\mathbf{m}} \int x^2 dx \prod_i \left( -2 \int_0^{\chi^*} d\chi_i W_i(\chi_i) \int \frac{k_i^2 dk_i}{(2\pi)^3} 4\pi (-1)^{l_i} j_{l_i}(k_i \chi_i) j_{l_i}(k_i x) \right) \\
&\quad \times B^\Psi(\{k_i\}, \{\eta_0 - \chi_i\}).
\end{aligned}$$

Now applying the Limber approximation again:

$$\begin{aligned}
\langle (\psi_1)_{l_1 m_1} (\psi_2)_{l_2 m_2} (\psi_3)_{l_3 m_3} \rangle &= (2\pi)^3 8 A_1^{\mathbf{m}} \int x^2 dx \prod_i \left( -2 \int_0^{\chi^*} d\chi_i W_i(\chi_i) \frac{1}{(2\pi)^3} \frac{\pi}{2\chi_i^2} \delta(x - \chi_i) 4\pi (-1)^{l_i} \right) \\
&\quad \times B^\Psi(\{l_i/\chi_i\}, \{\eta_0 - \chi_i\}) \\
&= (2\pi)^3 8 A_1^{\mathbf{m}} \int \chi^2 d\chi \prod_i \left( -2 W_i(\chi) \frac{1}{(2\pi)^3} \frac{\pi}{2\chi^2} 4\pi (-1)^{l_i} \right) B^\Psi(\{l_i/\chi\}, \eta_0 - \chi).
\end{aligned}$$

Finally, we aim to rewrite the above in terms of the angular bispectrum of the lensing potential.

The definition for the bispectrum of any set of randomly distributed spherical harmonic components  $X_{lm}^k$  is [8]

$$\langle (X_1)_{l_1 m_1} (X_2)_{l_2 m_2} (X_3)_{l_3 m_3} \rangle = \begin{pmatrix} l_1 & l_2 & l_3 \\ m_1 & m_2 & m_3 \end{pmatrix} B_{l_1 l_2 l_3}^{X_1 X_2 X_3}.$$

Note the independence on  $m_i$ , this necesarilly follows from statistical isotropy. If  $m_1 + m_2 + m_3 \neq 0$ , the associated Wigner-3j symbol vanishes and the bispectrum is set to zero. Also note that in this definition we have immediately generalized to include cross correlation between different fields  $X_1$ ,  $X_2$ ,  $X_3$ . This will be relevant later when we look at cross correlations between shear and convergence fields.

Using the above definition and the symmetry property

$$\begin{pmatrix} l_1 & l_2 & l_3 \\ -m_1 & -m_2 & -m_3 \end{pmatrix} = (-1)^{l_1+l_2+l_3} \begin{pmatrix} l_1 & l_2 & l_3 \\ m_1 & m_2 & m_3 \end{pmatrix}$$

we find

$$B_{l_1 l_2 l_3}^{\psi_1 \psi_2 \psi_3} = (-1)^{l_1 + l_2 + l_3} \sqrt{\frac{(2l_1 + 1)(2l_2 + 1)(2l_3 + 1)}{4\pi}} \begin{pmatrix} l_1 & l_2 & l_3 \\ 0 & 0 & 0 \end{pmatrix} (2\pi)^3 8 \\ \times \int \chi^2 d\chi \prod_i \left( -2W_i(\chi, \chi_*) \frac{1}{(2\pi)^3} \frac{\pi}{2\chi^2} 4\pi (-1)^{l_i} \right) B^\Psi(\{l_i/\chi\}, \{\eta_0 - \chi\}),$$

where we were able to drop the  $(-1)^{m_1 + m_2 + m_3}$  factor due to the bispectrum vanishing if that sum doesn't equal 0, as mentioned earlier. When all  $m_i$  equal zero, the Wigner 3j-symbol gains a number of useful properties. In particular, it vanishes if  $l_1 + l_2 + l_3$  is odd, meaning we can drop the  $(-1)^{l_1 + l_2 + l_3}$  factor. Additionally cancelling common factors then gives

$$B_{l_1 l_2 l_3}^{\psi_1 \psi_2 \psi_3} = -\sqrt{\frac{(2l_1 + 1)(2l_2 + 1)(2l_3 + 1)}{4\pi}} \begin{pmatrix} l_1 & l_2 & l_3 \\ 0 & 0 & 0 \end{pmatrix} 8 \int \frac{d\chi}{\chi^4} W_1(\chi) W_2(\chi) W_3(\chi) B^\Psi(\{l_i/\chi\}, \eta_0 - \chi). \quad (45)$$

### B.3 Gravitational potential spectra in terms of matter spectra

We can rewrite equations 44 and 45 in terms of the matter spectra instead of the gravitational potential spectra using the Poisson equation. This then allows us to numerically evaluate these objects using CAMB. The **density contrast** is defined as

$$\delta(\mathbf{x}) := \frac{\rho(\mathbf{x}) - \bar{\rho}}{\bar{\rho}}, \quad (46)$$

and the matter spectra are defined in terms of the Fourier transformed density contrast  $\delta(\mathbf{k})$  as

$$\langle \delta(\mathbf{k}, \eta) \delta(\mathbf{k}', \eta)^* \rangle = (2\pi)^3 \delta(\mathbf{k} - \mathbf{k}') P^\delta(\mathbf{k}, \eta), \\ \langle \delta(\mathbf{k}_1, \eta) \delta(\mathbf{k}_2, \eta) \delta(\mathbf{k}_3, \eta) \rangle = (2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) B^\delta(k_1, k_2, k_3, \eta).$$

The mean matter density of the universe,  $\bar{\rho}$ , can be expressed in terms of the matter density parameter  $\Omega_m$  and the critical density  $\rho_c$  as

$$\bar{\rho}(\eta) = \frac{3\Omega_m H_0^2}{8\pi G} \frac{1}{a(\eta)^3},$$

where  $a(\eta)$  is the only time dependent factor on the right hand side. In the case of a flat universe ( $K = 0$ ) we can thus rewrite the Poisson equation [10] as

$$\nabla^2 \Psi = 4\pi G a^2 \left( \frac{3\Omega_m H_0^2}{8\pi G} \frac{1}{a^3} \right) \delta = \frac{3\Omega_m H_0^2}{2} \frac{1}{a} \delta \implies \Psi(k, \eta) = -\frac{3\Omega_m H_0^2}{2} \frac{1}{a} \frac{\delta(k, \eta)}{k^2}, \quad (47)$$

where  $\Psi(k, \eta)$  and  $\delta(k, \eta)$  are functions in Fourier space. For the power- and bispectra we find

$$\langle \Psi(\mathbf{k}, \eta) \Psi^*(\mathbf{k}', \eta) \rangle = \frac{2\pi^2}{k^3} P^\Psi(k, \eta) \delta(\mathbf{k} - \mathbf{k}') \implies P^\Psi(k, \eta) = \frac{1}{k} (9\Omega_m^2 H_0^4 \pi) \frac{1}{a^2} P^\delta(k, \eta), \\ \langle \Psi(\mathbf{k}_1, \eta) \Psi(\mathbf{k}_2, \eta) \Psi(\mathbf{k}_3, \eta) \rangle = -(2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) B^\Psi(k_1, k_2, k_3, \eta) \\ \implies B^\Psi(k_1, k_2, k_3, \eta) = -\frac{1}{k_1^2 k_2^2 k_3^2} \left( \frac{3\Omega_m H_0^2}{2} \right)^3 \frac{1}{a^3} B^\delta(\{k_i\}, \eta).$$

Finally, we obtain:

$$\begin{aligned}
P_l^{\psi_X \psi_Y} &= \frac{9}{l^4} \Omega_m^2 H_0^4 \int_0^{\chi_*} \chi^2 d\chi a(\eta_0 - \chi)^{-2} W_X(\chi) W_Y(\chi) P^\delta(l/\chi, \eta_0 - \chi), \\
B_{l_1 l_2 l_3}^{\psi_X \psi_Y \psi_Z} &= \sqrt{\frac{(2l_1 + 1)(2l_2 + 1)(2l_3 + 1)}{4\pi}} \begin{pmatrix} l_1 & l_2 & l_3 \\ 0 & 0 & 0 \end{pmatrix} \frac{27}{l_1^2 l_2^2 l_3^2} \Omega_m^3 H_0^6 \\
&\quad \times \int \chi^2 d\chi a(\eta_0 - \chi)^{-3} W_X(\chi) W_Y(\chi) W_Z(\chi) B^\delta(\{l_i/\chi\}, \eta_0 - \chi).
\end{aligned}$$

## B.4 Convergence and shear multipole moments

**Convergence** Taking the harmonic decomposition of the convergence,

$$\kappa(\hat{\mathbf{n}}) = \kappa_{lm} Y_l^m(\hat{\mathbf{n}}), \quad (48)$$

you can relate it to the multipole moments of the lensing potential through

$$\kappa_{lm} = \frac{1}{2} (\nabla^2 \psi)_{lm} = \frac{l(l+1)}{2} \psi_{lm} \quad (49)$$

because the spherical harmonics are eigenfunctions of the Laplacian operator.

**Shear** On the other hand, the shear is given by

$$\gamma = \gamma_1 + i\gamma_2 = \frac{1}{2} \bar{\partial}_1 (\bar{\partial}_0 \psi) \quad (50)$$

where the spin raising operator,  $\bar{\partial}_s$  acts on a spin  $s$  function defined on  $S^2$  to create a spin  $s+1$  function. The above equality is proven explicitly in appendix ...  $\bar{\partial}_s$  can be written in  $(\theta, \phi)$  coordinates as

$$\bar{\partial}_s = -\sin^s \theta (\partial_\theta + \frac{i}{\sin \theta}) \frac{1}{\sin^s \theta}. \quad (51)$$

It is worth explaining that in this context a spin  $s$  function refers to a function  ${}_s f(\theta, \phi)$  that transforms under any rotation of coordinates by picking up a phase factor  $e^{is\alpha}$ , with  $\alpha$  the angle of the rotation, i.e.

$$f'(\theta', \phi') = e^{is\alpha} f(\theta, \phi). \quad (52)$$

Shear is actually a spin 2 function, which can immediately be seen by observing that rotating a galaxy image stretched and squeezed through weak lensing by 180 degrees gives the same stretching and squeezing, i.e. the same shear. It is thus expected that the shear is proportional to  $\bar{\partial}^2 \psi$ .

The spin weighted spherical harmonics,  ${}_s Y_l^m$  are eigenfunctions of the raising/lowering operator, in particular

$$\bar{\partial}_s Y_l^m = \sqrt{(l-s)(l+s+1)} {}_{s+1} Y_l^m. \quad (53)$$

Decomposing the shear as

$$\gamma(\theta, \phi) = \gamma_{lm} Y_l^m(\theta, \phi), \quad (54)$$

we thus find

$$\gamma_{lm} = \frac{1}{2} \sqrt{(l-1)(l)(l+1)(l+2)} \psi_{lm} \quad (55)$$

## C Fisher Matrix Analysis

### C.1 Determining uncertainty in experimental parameters

Fisher matrix formalism is used to find a lower bound on the constraints we can place on experimental parameters. It combines the Cramer-Rao Inequality [7] with the assumption that we have unbiased estimators following a gaussian distribution [10]. In particular, denoting the Fisher matrix by  $F$  and the parameters as  $\theta_i$ , it can be shown that

$$\text{Var}_{\theta_i}(\hat{\theta}_i) \geq (F^{-1})_{\theta_i\theta_j}. \quad (56)$$

In the case of  $n$  potentially different random variables  $x_i$ , each with associated mean  $\mu_{x_i}$  and error  $\sigma_{x_i}$ , the Fisher matrix is given as

$$F_{\theta_i\theta_j} := \sum_{p,q=1}^n \frac{\partial \mu_{x_p}}{\partial \theta_i}(\tilde{\theta}_k) (\text{Cov}^{-1})_{x_px_q} \frac{\partial \mu_{x_q}}{\partial \theta_j}(\tilde{\theta}_k), \quad (57)$$

where Cov is the covariance matrix associated with the random vector  $(x_1, \dots, x_n)$ ,  $\text{Cov}_{x_px_q} := \text{Cov}(x_p, x_q)$ .

### C.2 Fisher matrices for power- bispectra with multiple tracers

When considering power spectra, the definition of the Fisher matrix immediately gives

$$F_{\alpha\beta} = \sum_{l_{\min} \leq l, l' \leq l_{\max}} \sum_{[XY][X'Y']} \partial_\alpha C_l^{XY} (\text{Cov}^{-1})_{l,l'}^{XY,X'Y'} \partial_\beta C_{l'}^{X'Y'}.$$

Here the covariance matrix is given as

$$\text{Cov}_{l,l'}^{XY,X'Y'} = \langle \hat{C}_l^{XY} \hat{C}_{l'}^{X'Y'} \rangle.$$

The estimator of a power or bispectrum is just the product of the estimators of the appropriate tracers,  $\hat{C}_l^{XY} = \hat{X}(l)\hat{Y}(l)$ . The sum over  $[XY]$  and  $[X'Y']$  denotes a sum over possible tracer combinations *without* counting permutations of tracer configuration. This is because permutations are not distinct signals, i.e.  $\hat{X}(l)\hat{Y}(l) = \hat{Y}(l)\hat{X}(l)$ . In fact, if we were to count these permutations as distinct signals we will get identical columns and/or rows in our covariance matrix making inversion impossible:

$$\langle \hat{X}(l)\hat{Y}(l)\hat{X}'(l')\hat{Y}'(l') \rangle = \langle \hat{Y}(l)\hat{X}(l)\hat{X}'(l')\hat{Y}'(l') \rangle, \quad \forall X', Y', l'.$$

To evaluate the covariance matrix we again assume that the estimators are gaussian distributed so that we can do apply a wick contraction, as is commonly done [25]. In this case we get

$$\text{Cov}_{l,l'}^{XY,X'Y'} = \delta_{ll'} \left( C_l^{XX'} C_l^{YY'} + C_l^{XY'} C_l^{YX'} \right).$$

Under the Gaussian approximation the gaussian matrix thus vanishes except for  $3 \times 3$  block matrices (in the case of 2 tracers) on the diagonal. The fisher matrix is then

$$F_{\alpha\beta} = \sum_l \sum_{[XY][X'Y']} \partial_\alpha C_l^{XY} (\text{Cov}_l^{-1})^{XY,X'Y'} \partial_\beta C_l^{X'Y'}.$$

$\text{Cov}_l^{-1}$  here denotes the inverse of the block matrix at  $l$ .

Next, we consider the Fisher matrix for bispectra measurements.

$$F_{\alpha\beta} = \sum_{\text{distinct signals}} \sum_{\text{distinct signals prime}} B_{l_1 l_2 l_3}^{XYZ} (\text{Cov}^{-1})_{l_1 l_2 l_3, l'_1 l'_2 l'_3}^{XYZ, X' Y' Z'} B_{l'_1 l'_2 l'_3}^{X' Y' Z'}.$$

Counting only distinct signals requires more care compared to the power spectra. The rule is that  $B_{l_1 l_2 l_3}^{XYZ}$  and  $B_{l'_1 l'_2 l'_3}^{X' Y' Z'}$  are not distinct signals if there exists a permutation  $\sigma$  that simultaneously maps  $X' Y' Z'$  to  $XYZ$  and  $l'_1 l'_2 l'_3$  to  $l_1 l_2 l_3$ . It turns out that we can write a sum over distinct signals out as

$$\sum_{\text{distinct signals}} = \underbrace{\sum_{l_1=l_2=l_3} \sum_{[XYZ]}}_{\text{sum 1}} + \underbrace{\sum_{l_1=l_2 \neq l_3} \sum_{[XY]Z}}_{\text{sum 2}} + \underbrace{\sum_{l_1 < l_2 < l_3} \sum_{XYZ}}_{\text{sum 3}}.$$

With the same definition again for the  $[\cdot]$  notation. For example:

$$\{[XY]Z | X, Y, Z \in \{\psi_\kappa, \psi_\gamma\}\} = \{\psi_\kappa \psi_\kappa \psi_\kappa, \psi_\kappa \psi_\gamma \psi_\kappa, \psi_\gamma \psi_\gamma \psi_\kappa, \psi_\kappa \psi_\kappa \psi_\gamma, \psi_\kappa \psi_\gamma \psi_\gamma, \psi_\gamma \psi_\gamma \psi_\gamma\}.$$

It follows to show that the sets that these sums sum over form a partition of the set of all distinct signals. Clearly all 3 sets are pairwise disjoint (no common elements) because of the criteria for the  $l_i$ 's. To show that their union covers the set of distinct signals, start by considering an arbitrary signal. Its  $l$  configuration will trivially correspond to exactly one of the three sums. If it is sum 1, then we are free to permute the  $XYZ$ 's by virtue of the  $l$ 's being identical so we will be able to match the  $XYZ$  configuration to one of the elements of  $\{\{XYZ\}\}$ . Similarly, if the  $l$  configuration corresponds to sum 2, then we are free to permute the  $XY$  configuration to match with one of the elements of  $\{\{XY\}Z\}$ . The  $Z$  value does not matter because any  $Z$  value is accounted for. For sum 3 we can argue that we can switch around the order of the  $l$ 's to satisfy  $l_1 < l_2 < l_3$  and the corresponding  $XYZ$  configuration will be accounted for in sum 3 because all  $XYZ$  combinations are accounted. Finally, it is simple to verify that no distinct signal is counted more than once within each sum.

**Covariance matrix** Every element of the Fisher matrix can be seen as an inner product weighted by the inverse covariance matrix. We can choose how we order the vectors<sup>3</sup>. We organize the vectors according to the sum 1, 2, and 3 parts first. Then by  $l$  configuration. Within each  $l$  configuration we can choose any ordering for the  $XYZ$  configurations. The covariance matrix now becomes a block matrix with each block corresponding to a  $l_i$  and  $l'_i$  configuration. When wick contracting using the gaussian approximation, every block matrix where  $l_1 l_2 l_3$  is not a permutation of  $l'_1 l'_2 l'_3$  vanishes. The entries of each block matrix are given as

$$\begin{aligned} (\text{Cov}_{l_1 l_2 l_3})^{XYZ, X' Y' Z'} &= C_{l_1}^{XX'} C_{l_2}^{YY'} C_{l_3}^{ZZ'} + \delta_{l_1 l_2} C_{l_1}^{XY'} C_{l_2}^{YX'} C_{l_3}^{ZZ'} + \delta_{l_2 l_3} C_{l_1}^{XX'} C_{l_2}^{YZ'} C_{l_3}^{ZX'} \\ &+ \delta_{l_3 l_1} C_{l_1}^{XZ'} C_{l_2}^{ZY'} C_{l_3}^{ZX'} + \delta_{l_1 l_2} \delta_{l_2 l_3} \left( C_{l_1}^{XY'} C_{l_2}^{YZ'} C_{l_3}^{ZX'} + C_{l_1}^{XZ'} C_{l_2}^{YX'} C_{l_3}^{ZY'} \right). \end{aligned}$$

With our ordering this means that the covariance matrix is again a diagonal block matrix, now with blocks of size  $4 \times 4$  (sum 1),  $6 \times 6$  (sum 2), and  $8 \times 8$  (sum 3).

<sup>3</sup>The entries are the derivatives of the bispectra

### C.3 Explicit form for inverse covariance matrix

The Fisher matrix above can be significantly simplified and written as

$$F_{\alpha\beta} = \sum_{l_1 \leq l_2 \leq l_3} \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \sum_{XYZ} \sum_{X'Y'Z'} \partial_\alpha B_{l_1 l_2 l_3}^{XYZ} (C^{-1})_{l_1}^{XX'} (C^{-1})_{l_2}^{YY'} (C^{-1})_{l_3}^{ZZ'} \partial_\beta B_{l_1 l_2 l_3}^{X'Y'Z'}$$

where

$$C_l := \begin{pmatrix} C_l^{\psi_\kappa \psi_\kappa} & C_l^{\psi_\kappa \psi_\gamma} \\ C_l^{\psi_\kappa \psi_\gamma} & C_l^{\psi_\gamma \psi_\gamma} \end{pmatrix}$$

and  $\mathcal{P}_{l_1 l_2 l_3}$  is defined as the number of distinct permutations that can be made with  $l_1 l_2 l_3$ . This form was, for example, used in [17]<sup>4</sup>.

To show that the above is the same as the formula for the Fisher matrix given earlier, first partition the sum in the same way and collect all terms that fit in the different categories.

$$\begin{aligned} F_{\alpha\beta} &= \sum_{l_1=l_2=l_3} \sum_{[XYZ]} \sum_{[X'Y'Z']} \partial_\alpha B_{l_1 l_2 l_3}^{XYZ} \left( \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \sum_{d.p.XYZ} \sum_{d.p.X'Y'Z'} (C^{-1})_{l_1}^{XX'} (C^{-1})_{l_2}^{YY'} (C^{-1})_{l_3}^{ZZ'} \right) \partial_\beta B_{l_1 l_2 l_3}^{X'Y'Z'} \\ &+ \sum_{l_1=l_2 \neq l_3} \sum_{[XY]Z} \sum_{[X'Y']Z'} \partial_\alpha B_{l_1 l_2 l_3}^{XYZ} \left( \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \sum_{d.p.XY} \sum_{d.p.X'Y'} (C^{-1})_{l_1}^{XX'} (C^{-1})_{l_2}^{YY'} (C^{-1})_{l_3}^{ZZ'} \right) \partial_\beta B_{l_1 l_2 l_3}^{X'Y'Z'} \\ &+ \sum_{l_1 < l_2 < l_3} \sum_{XYZ} \sum_{X'Y'Z'} \partial_\alpha B_{l_1 l_2 l_3}^{XYZ} (C^{-1})_{l_1}^{XX'} (C^{-1})_{l_2}^{YY'} (C^{-1})_{l_3}^{ZZ'} \partial_\beta B_{l_1 l_2 l_3}^{X'Y'Z'}, \end{aligned}$$

where “d.p.” stands for “distinct permutations”. The entries above are then the entries of the inverses of the block matrices mentioned earlier. This can be checked. For example, for the  $l_1 = l_2 = l_3$  sum the multiplication of block matrices corresponding to the same  $l_i$  configuration can be written as:

$$\begin{aligned} &\sum_{[X'Y'Z']} \left( \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \sum_{d.p.XYZ} \sum_{d.p.X'Y'Z'} (C^{-1})_l^{XX'} (C^{-1})_l^{YY'} (C^{-1})_l^{ZZ'} \right) (C_l^{X'X''} C_l^{Y'Y''} C_l^{Z'Z''} + \text{perms } X''Y''Z'') \\ &= \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \left[ \left( \sum_{d.p.XYZ} \sum_{X'Y'Z'} (C^{-1})_l^{XX'} (C^{-1})_l^{YY'} (C^{-1})_l^{ZZ'} \right) C_l^{X'X''} C_l^{Y'Y''} C_l^{Z'Z''} \right] + \text{perms } X''Y''Z'' \\ &= \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \sum_{d.p.XYZ} \delta_{XX''} \delta_{YY''} \delta_{ZZ''} + \text{perms } X''Y''Z'' = \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} \delta_{[XYZ], [X''Y''Z'']} + \text{perms } X''Y''Z'' \\ &= \delta_{[XYZ], [X''Y''Z'']}. \end{aligned}$$

The sum over the different wick contractions will similarly cancel with the  $\mathcal{P}_{l_1 l_2 l_3}/6$  factor for the  $l_1 = l_2 \neq l_3$  sum. For the  $l_1 < l_2 < l_3$  sum the proof is similar as well except no cancellation is required.

The same type of simplification can be made in the Fisher matrix for the power spectrum, though it doesn't offer any significant benefits compared to our current  $3 \times 3$  block matrix form.

<sup>4</sup>Note that in [17] this is based on a previous equation summing over *all*  $l_i$  (so including permutations of each configuration) which is missing a factor of  $1/6$ .

## C.4 Signal to Noise Ratio (SNR)

To measure the overall detectability of the lensing spectra, we introduce an overall amplitude of our signal,  $A$ , with fiducial value 1 as experimental parameter and compute  $F_{AA}$ . Obviously,  $\partial_A (AB_{l_1 l_2 l_3}^{XYZ})|_{A=1} = B_{l_1 l_2 l_3}^{XYZ}$ , so we find

$$\left(\frac{S}{N}\right)^2 := F_{AA} = \sum_{XYZ, X'Y'Z'} \sum_{l_1 \leq l_2 \leq l_3} \frac{\mathcal{P}_{l_1 l_2 l_3}}{6} B_{l_1 l_2 l_3}^{XYZ} (C^{-1})_{l_1}^{XX'} (C^{-1})_{l_2}^{YY'} (C^{-1})_{l_3}^{ZZ'} B_{l_1 l_2 l_3}^{X'Y'Z'}.$$

The equation for the SNR of the power spectra is identical in form.

## C.5 Combined Fisher Matrix

To compute the Fisher matrix of an experiment measuring both lensing power- and bispectra we are required to compute and invert the full covariance matrix. Under the Gaussian matrix this simplifies trivially. The correlation between the a power- and bispectrum estimator contains an odd (5) amount of fields and thus always vanishes. We are thus allowed to simply add the Fisher matrices of the power- and bispectra to compute the combined Fisher matrix.

## C.6 Fisher Matrices and Eigenvalues

As explained earlier, inverting the Fisher matrix gives us a covariance matrix of the estimators,  $\hat{\theta}_i$  of our experimental parameters. For any vector  $\mathbf{v}$ ,

$$\mathbf{v}_j^T \text{Cov}_{ij} \mathbf{v}_i = \langle (v_i \hat{\theta}_i)^2 \rangle.$$

In particular, if the covariance matrix thus has some large eigenvalue  $\lambda_L$  (large compared to the overall accuracy of our experiment), with eigenvector  $v_i$ , then the corresponding estimator,  $v_i \hat{\theta}_i$ , has a large variance and thus represents an approximate degeneracy in our parameters. Clearly, when looking at the Fisher matrix we thus instead look for the smallest eigenvalue.

If we have such an abnormally small eigenvalue we are likely to run into problems with calculating and inverting the Fisher matrix due to numerical errors. These problems are that inversion becomes unstable and that the matrix as a whole might not be PSD because the small eigenvalue was (wrongly) calculated to be negative. These problems are likely because (1) the smallest eigenvalue is sensitive to numerical errors due to being small and (2) the errors on the eigenvalues is proportional to the condition number of the matrix [13], where the condition number is the ratio between the largest and smallest eigenvalue of the matrix [26]. Specifically, for a matrix  $A$  with some perturbation  $\Delta A$ , the corresponding perturbation of eigenvalue  $\lambda_i$  is bounded above as

$$|\Delta \lambda_i| \leq \|\Delta A\| \kappa(A),$$

with  $\kappa(A)$  the condition number.

Before inverting a Fisher matrix, it is thus recommended to remove rows and columns corresponding to parameters that are poorly constrained by the experiment, which can be determined by checking for small eigenvalues and their eigenvectors.

## D Shear equals twice spin raised lensing potential

Consider a point on  $S^2$ ,  $(r_0, \theta_0, \phi_0)$ , at which we observe some cosmological object. We can then define a set of cartesian coordinates  $(\tilde{r}, \tilde{y}, \tilde{x})$  as shown in figure 5.

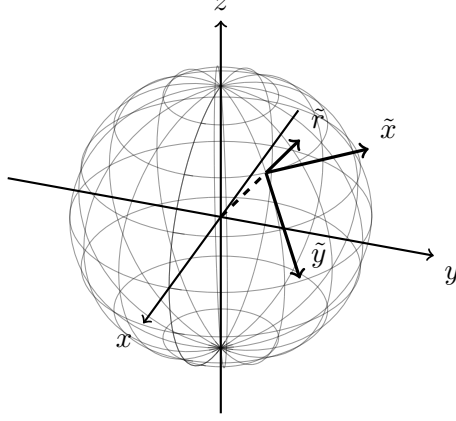


Figure 5:  $(\tilde{r}, \tilde{y}, \tilde{x})$  coordinates defined for a point on the unit sphere. These act as ordinary cartesian coordinates but rotated such that, at the associated point on  $S^2$ ,  $\hat{\tilde{r}}$  points straight out of the unit sphere,  $\hat{\tilde{y}}$  is parallel to the great arc with constant  $\phi$  and  $\hat{\tilde{x}}$  is parallel to the great arc with constant  $\theta$ . These coordinates are used to define the shear and convergence in terms of the lensing potential.

Note that there it isn't obvious whether to define these coordinates at the point where the lensed light hits  $S^2$  or the unlensed light hits  $S^2$ . We will assume that lensing effects are sufficiently weak that either definition works. We can then express the tilde coordinates in terms of spherical coordinates either by applying a rotation matrix or by calculating the  $r, \theta, \phi$  derivatives of  $(x, y, z)$  coordinates at  $(r_0, \theta_0, \phi_0)$  to find  $\hat{\tilde{r}}$ ,  $\hat{\tilde{\theta}}$ , and  $\hat{\tilde{\phi}}$  and then take inner products. Regardless, we find

$$\tilde{r} = r \sin \theta \cos \phi \sin \theta_0 \cos \phi_0 + r \sin \theta \sin \phi \sin \theta_0 \sin \phi_0 + r \cos \theta \cos \theta_0, \quad (58)$$

$$\tilde{y} = r \sin \theta \cos \phi \cos \theta_0 \cos \phi_0 + r \sin \theta \sin \phi \cos \theta_0 \sin \phi_0 - r \cos \theta \sin \theta_0, \quad (59)$$

$$\tilde{x} = -r \sin \theta \cos \phi \sin \phi_0 + r \sin \theta \sin \phi \cos \phi_0. \quad (60)$$

This gives derivatives

$$\begin{aligned} \frac{\partial}{\partial \theta} &= (r \cos \theta \cos \phi \sin \theta_0 \cos \phi_0 + r \cos \theta \sin \phi \sin \theta_0 \sin \phi_0 - r \sin \theta \cos \theta_0) \frac{\partial}{\partial \tilde{r}} \\ &\quad + (r \cos \theta \cos \phi \cos \theta_0 \cos \phi_0 + r \cos \theta \sin \phi \cos \theta_0 \sin \phi_0 + r \sin \theta \sin \theta_0) \frac{\partial}{\partial \tilde{y}} \\ &\quad + (-r \cos \theta \cos \phi \sin \phi_0 + r \cos \theta \sin \phi \cos \phi_0) \frac{\partial}{\partial \tilde{x}}. \\ \frac{\partial}{\partial \phi} &= (-r \sin \theta \sin \phi \sin \theta_0 \cos \phi_0 + r \sin \theta \cos \phi \sin \theta_0 \sin \phi_0) \frac{\partial}{\partial \tilde{r}} \\ &\quad + (-r \sin \theta \sin \phi \cos \theta_0 \cos \phi_0 + r \sin \theta \cos \phi \cos \theta_0 \sin \phi_0) \frac{\partial}{\partial \tilde{y}} \\ &\quad + (r \sin \theta \sin \phi \sin \phi_0 + r \sin \theta \cos \phi \cos \phi_0) \frac{\partial}{\partial \tilde{x}}. \end{aligned}$$



Evaluated at our point of interest we obtain  $\partial_\theta = \partial_{\tilde{y}}$  and  $\partial_\phi = \sin \theta_0 \partial_{\tilde{x}}$ . The second order derivatives can then be obtained using the first order derivatives. We can immediately evaluate them at the point to get

$$\begin{aligned}\partial_\phi^2|_{(r_0, \theta_0, \phi_0)} &= -\sin^2 \theta_0 \partial_{\tilde{r}} - \sin \theta_0 \cos \theta_0 \partial_{\tilde{y}} + \sin^2 \theta_0 \partial_{\tilde{x}}^2, \\ \partial_\theta \partial_\phi|_{(r_0, \theta_0, \phi_0)} &= \cos \theta_0 \partial_{\tilde{x}} + \sin \theta_0 \partial_{\tilde{x}} \partial_{\tilde{y}}, \\ \partial_\theta^2|_{(r_0, \theta_0, \phi_0)} &= -\partial_{\tilde{r}} + \partial_{\tilde{y}}^2.\end{aligned}$$

Thus, at  $(r_0, \theta_0, \phi_0)$ ,

$$\begin{aligned}\frac{1}{2} \tilde{\partial}_1(\tilde{\partial}_0 \psi) &= \frac{1}{2} \sin \theta (\partial_\theta + \frac{i}{\sin \theta} \partial_\phi) (\frac{1}{\sin \theta} (\partial_\theta + \frac{1}{\sin \theta} \partial_\phi)) \\ &= \frac{\partial^2 \psi}{\partial \theta^2} - \frac{\cos \theta}{\sin \theta} \frac{\partial \psi}{\partial \theta} + \frac{2i}{\sin \theta} \frac{\partial^2 \psi}{\partial \theta \partial \phi} - \frac{1}{\sin^2 \theta} \frac{\partial^2 \psi}{\partial \phi^2} - 2i \frac{\cos \theta}{\sin^2 \theta} \frac{\partial \psi}{\partial \phi} \\ &= \frac{1}{2} (\partial_{\tilde{y}}^2 - \partial_{\tilde{x}}^2 + 2i \partial_{\tilde{x}} \partial_{\tilde{y}}) \psi = \gamma_1 + i \gamma_2 = \gamma,\end{aligned}$$

as required.