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Consistency relations for large-scale structures: applications to the integrated Sachs-Wolfe effect and the kinematic Sunyaev-Zeldovich effect

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ABSTRACT

Consistency relations of large-scale structures provide exact nonperturbative results for cross-correlations of cosmic fields in the squeezed limit. They only depend on the equivalence principle and the assumption of Gaussian initial conditions, and remain nonzero at equal times for cross-correlations of density fields with velocity or momentum fields, or with the time derivative of density fields. We show how to apply these relations to observational probes that involve the integrated Sachs-Wolfe effect or the kinematic Sunyaev-Zeldovich effect. In the squeezed limit, this allows us to express the three-point cross-correlations, or bispectra, of two galaxy or matter density fields, or weak lensing convergence fields, with the secondary CMB distortion in terms of products of a linear and a nonlinear power spectrum. In particular, we find that cross-correlations with the integrated Sachs-Wolfe effect show a specific angular dependence. These results could be used to test the equivalence principle and the primordial Gaussianity, or to check the modeling of large-scale structures.

Key words. Cosmology – large-scale structure of the Universe

1. Introduction

Measuring statistical properties of cosmological structures is not only an efficient tool to describe and understand the main components of our Universe, but also it is a powerful probe of possible new physics beyond the standard Λ CDM concordance model. However, on large scales cosmological structures are described by perturbative methods, while smaller scales are described by phenomenological models or studied with numerical simulations. It is therefore difficult to obtain accurate predictions on the full range of scales probed by galaxy and lensing surveys. Furthermore, if we consider galaxy density fields, theoretical predictions remain sensitive to the galaxy bias which involves phenomenological modeling of star formation, even if we use cosmological numerical simulations. As a consequence, exact analytical results that go beyond low-order perturbation theory and also apply to biased tracers are very rare.

Recently, some exact results have been obtained (Kehagias & Riotto 2013; Peloso & Pietroni 2013; Creminelli et al. 2013; Kehagias et al. 2014a; Peloso & Pietroni 2014; Creminelli et al. 2014; Valageas 2014b; Horn et al. 2014, 2015) in the form of “kinematic consistency relations”. They relate the $(\ell + n)$ -density correlation, with ℓ large-scale wave numbers and n small-scale wave numbers, to the n -point small-scale density correlation. These relations, obtained at the leading order over the large-scale wave numbers, arise from the Equivalence Principle (EP) and the assumption of Gaussian initial conditions. The equivalence principle ensures that small-scale structures respond to a large-scale perturbation by a uniform displacement while primordial Gaussianity provides a simple relation between correlation and response functions (see Valageas et al. (2016) for the additional terms associated with

non-Gaussian initial conditions). Hence, such relations express a kinematic effect that vanishes for equal-times statistics, as a uniform displacement has no impact on the statistical properties of the density field observed at a given time.

In practice, it is however difficult to measure different-times density correlations and it would therefore be useful to obtain relations that remain nonzero at equal times. One possibility to overcome such problem, is to go to higher orders and take into account tidal effects, which at leading order are given by the response of small-scale structures to a change of the background density. Such an approach, however, introduces some additional approximations (Valageas 2014a; Kehagias et al. 2014b; Nishimichi & Valageas 2014).

Fortunately, it was recently noticed that by cross-correlating density fields with velocity or momentum fields, or with the time derivative of the density field, one obtains consistency relations that do not vanish at equal times (Rizzo et al. 2016). Indeed, the kinematic effect modifies the amplitude of the large-scale velocity and momentum fields, while the time derivative of the density field is obviously sensitive to different-times effects.

In this paper, we investigate the observational applicability of these new relations. We consider the lowest-order relations, which relate three-point cross-correlations or bispectra in the squeezed limit to products of a linear and a nonlinear power spectrum. To involve the non-vanishing consistency relations, we study two observable quantities, the secondary anisotropy Δ_{ISW} of the cosmic microwave background (CMB) radiation due to the integrated Sachs-Wolfe effect (ISW), and the secondary anisotropy Δ_{kSZ} due to the kinematic Sunyaev-Zeldovich (kSZ) effect. The first process, associated with the motion of CMB photons through time-dependent gravitational potentials, depends on the time derivative of the matter density field. The second pro-

cess, associated with the scattering of CMB photons by free electrons, depends on the free electrons velocity field. We investigate the cross correlations of these two secondary anisotropies with both galaxy density fields and the cosmic weak lensing convergence.

This paper is organized as follows. In section 2 we recall the consistency relations of large-scale structures that apply to density, momentum and momentum-divergence (i.e., time derivative of the density) fields. We describe the various observational probes that we consider in this paper in section 3. We study the ISW effect in section 4 and the kSZ effect in section 5. We conclude in section 6.

2. Consistency relations for large-scale structures

2.1. Consistency relations for density correlations

As described in recent works (Kehagias & Riotto 2013; Peloso & Pietroni 2013; Creminelli et al. 2013; Kehagias et al. 2014a; Peloso & Pietroni 2014; Creminelli et al. 2014; Valageas 2014b; Horn et al. 2014, 2015), it is possible to obtain exact relations between density correlations of different orders in the squeezed limit, where some of the wavenumbers are in the linear regime and far below the other modes that may be strongly nonlinear. These “kinematic consistency relations”, obtained at the leading order over the large-scale wavenumbers, arise from the equivalence principle and the assumption of Gaussian primordial perturbations. They express the fact that, at leading order where a large-scale perturbation corresponds to a linear gravitational potential (hence a constant Newtonian force) over the extent of a small-size structure, the latter falls without distortions in this large-scale potential.

Then, in the squeezed limit $k \rightarrow 0$, the correlation between one large-scale density mode $\tilde{\delta}(\mathbf{k})$ and n small-scale density modes $\tilde{\delta}(\mathbf{k}_j)$ can be expressed in terms of the n -point small-scale correlation, as

$$\begin{aligned} \langle \tilde{\delta}(\mathbf{k}, \eta) \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \rangle'_{k \rightarrow 0} &= -P_L(k, \eta) \langle \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \rangle' \\ &\times \sum_{i=1}^n \frac{D(\eta_i)}{D(\eta)} \frac{\mathbf{k}_i \cdot \mathbf{k}}{k^2}, \end{aligned} \quad (1)$$

where the tilde denotes the Fourier transform of the fields, η is the conformal time, $D(\eta)$ is the linear growth factor, the prime in $\langle \dots \rangle'$ denotes that we factored out the Dirac factor, $\langle \dots \rangle = \langle \dots \rangle' \delta_D(\sum \mathbf{k}_j)$, and $P_L(k)$ is the linear matter power spectrum. It is worth stressing that these relations are valid even in the nonlinear regime and for biased galaxy fields $\tilde{\delta}_g(\mathbf{k}_j)$. The right-hand side gives the squeezed limit of the $(1+n)$ correlation at the leading order, which scales as $1/k$. It vanishes at this order at equal times, because of the constraint associated with the Dirac factor $\delta_D(\sum \mathbf{k}_j)$.

The geometrical factors $(\mathbf{k}_i \cdot \mathbf{k})$ vanish if $\mathbf{k}_i \perp \mathbf{k}$. Indeed, the large-scale mode induces a uniform displacement along the direction of \mathbf{k} . This has no effect on small-scale plane waves of wavenumbers \mathbf{k}_i with $\mathbf{k}_i \perp \mathbf{k}$, as they remain identical after such a displacement. Therefore, the terms in the right-hand side of Eq.(1) must vanish in such orthogonal configurations, as we can check from the explicit expression.

The simplest relation that one can obtain from Eq.(1) is for the bispectrum with $n = 2$,

$$\begin{aligned} \langle \tilde{\delta}(\mathbf{k}, \eta) \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\delta}_g(\mathbf{k}_2, \eta_2) \rangle'_{k \rightarrow 0} &= -P_L(k, \eta) \frac{\mathbf{k}_1 \cdot \mathbf{k}}{k^2} \\ &\times \langle \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\delta}_g(\mathbf{k}_2, \eta_2) \rangle' \frac{D(\eta_1) - D(\eta_2)}{D(\eta)}, \end{aligned} \quad (2)$$

where we used that $\mathbf{k}_2 = -\mathbf{k}_1 - \mathbf{k} \rightarrow -\mathbf{k}_1$. For generality, we considered here the small-scale fields $\tilde{\delta}_g(\mathbf{k}_1)$ and $\tilde{\delta}_g(\mathbf{k}_2)$ to be associated with biased tracers such as galaxies. The tracers associated with \mathbf{k}_1 and \mathbf{k}_2 can be different and have different bias. At equal times the right-hand side of Eq.(2) vanishes, as recalled above.

2.2. Consistency relations for momentum correlations

The density consistency relations (1) express the uniform motion of small-scale structures by large-scale modes. This simple kinematic effect vanishes for equal-time correlations of the density field, precisely because there are no distortions, while there is a nonzero effect at different times because of the motion of the small-scale structure between different times. However, as pointed out in Rizzo et al. (2016), it is possible to obtain non-trivial equal-times results by considering velocity or momentum fields, which are not only displaced but also see their amplitude affected by the large-scale mode. Let us consider the momentum \mathbf{p} defined by

$$\mathbf{p} = (1 + \delta)\mathbf{v}, \quad (3)$$

where \mathbf{v} the peculiar velocity. Then, in the squeezed limit $k \rightarrow 0$, the correlation between one large-scale density mode $\tilde{\delta}(\mathbf{k})$, n small-scale density modes $\tilde{\delta}(\mathbf{k}_j)$, and m small-scale momentum modes $\tilde{\mathbf{p}}(\mathbf{k}_j)$ can be expressed in terms of $(n+m)$ small-scale correlations, as

$$\begin{aligned} \langle \tilde{\delta}(\mathbf{k}, \eta) \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \prod_{j=n+1}^{n+m} \tilde{\mathbf{p}}(\mathbf{k}_j, \eta_j) \rangle'_{k \rightarrow 0} &= -P_L(k, \eta) \\ &\times \left\{ \langle \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \prod_{j=n+1}^{n+m} \tilde{\mathbf{p}}(\mathbf{k}_j, \eta_j) \rangle' \sum_{i=1}^{n+m} \frac{D(\eta_i)}{D(\eta)} \frac{\mathbf{k}_i \cdot \mathbf{k}}{k^2} \right. \\ &+ \sum_{i=n+1}^{n+m} \frac{(dD/dn)(\eta_i)}{D(\eta)} \langle \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \prod_{j=n+1}^{i-1} \tilde{\mathbf{p}}(\mathbf{k}_j, \eta_j) \\ &\left. \times \left(i \frac{\mathbf{k}}{k^2} [\delta_D(\mathbf{k}_i) + \tilde{\delta}(\mathbf{k}_i, \eta_i)] \prod_{j=i+1}^{n+m} \tilde{\mathbf{p}}(\mathbf{k}_j, \eta_j) \right) \right\}. \end{aligned} \quad (4)$$

These relations are again valid in the nonlinear regime and for biased galaxy fields $\tilde{\delta}_g(\mathbf{k}_j)$ and $\tilde{\mathbf{p}}_g(\mathbf{k}_j)$. As for the density consistency relation (1), the first term vanishes at this order at equal times. The second term however, which arises from the $\tilde{\mathbf{p}}$ fields only, remains nonzero. This is due to the fact that $\tilde{\mathbf{p}}$ involves the velocity, the amplitude of which is affected by the motion induced by the large-scale mode.

The simplest relation associated with Eq.(4) is the bispectrum among two density-contrast fields and one momentum field,

$$\begin{aligned} \langle \tilde{\delta}(\mathbf{k}, \eta) \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\mathbf{p}}_g(\mathbf{k}_2, \eta_2) \rangle'_{k \rightarrow 0} &= -P_L(k, \eta) \\ &\times \left(\frac{\mathbf{k}_1 \cdot \mathbf{k}}{k^2} \langle \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\mathbf{p}}_g(\mathbf{k}_2, \eta_2) \rangle' \frac{D(\eta_1) - D(\eta_2)}{D(\eta)} + \right. \\ &\left. + i \frac{\mathbf{k}}{k^2} \langle \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\delta}_g(\mathbf{k}_2, \eta_2) \rangle' \frac{1}{D(\eta)} \frac{dD}{d\eta}(\eta_2) \right). \end{aligned} \quad (5)$$

For generality, we considered here the small-scale fields $\tilde{\delta}_g(\mathbf{k}_1)$ and $\tilde{\mathbf{p}}_g(\mathbf{k}_2)$ to be associated with biased tracers such as galaxies, and the tracers associated with \mathbf{k}_1 and \mathbf{k}_2 can again be different and have different bias. At equal times Eq.(5) reads as

$$\langle \tilde{\delta}(\mathbf{k}) \tilde{\delta}_g(\mathbf{k}_1) \tilde{\mathbf{p}}_g(\mathbf{k}_2) \rangle'_{k \rightarrow 0} = -i \frac{\mathbf{k}}{k^2} \frac{d \ln D}{d\eta} P_L(k) P_g(k_1), \quad (6)$$

where $P_g(k)$ is the galaxy nonlinear power spectrum and we omitted the common time dependence. This result does not vanish thanks to the term generated by $\tilde{\mathbf{p}}$ in the consistency relation (5).

2.3. Consistency relations for momentum-divergence correlations

In addition to the momentum field \mathbf{p} , we can consider its divergence λ , defined by

$$\lambda \equiv \nabla \cdot [(1 + \delta)\mathbf{v}] = -\frac{\partial \delta}{\partial \eta}. \quad (7)$$

The second equality expresses the continuity equation, that is, the conservation of matter. In the squeezed limit we obtain from Eq.(4) (Rizzo et al. 2016)

$$\begin{aligned} \langle \tilde{\delta}(\mathbf{k}, \eta) \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \prod_{j=n+1}^{n+m} \tilde{\lambda}(\mathbf{k}_j, \eta_j) \rangle'_{k \rightarrow 0} &= -P_L(k, \eta) \\ &\times \left\{ \left\langle \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \prod_{j=n+1}^{n+m} \tilde{\lambda}(\mathbf{k}_j, \eta_j) \right\rangle' \sum_{i=1}^{n+m} \frac{D(\eta_i)}{D(\eta)} \frac{\mathbf{k}_i \cdot \mathbf{k}}{k^2} \right. \\ &\quad - \sum_{i=n+1}^{n+m} \langle \tilde{\delta}(\mathbf{k}_i, \eta_i) \prod_{j=1}^n \tilde{\delta}(\mathbf{k}_j, \eta_j) \prod_{\substack{j=n+1 \\ j \neq i}}^{n+m} \tilde{\lambda}(\mathbf{k}_j, \eta_j) \rangle' \\ &\quad \left. \times \frac{(dD/d\eta)(\eta_i)}{D(\eta)} \frac{\mathbf{k}_i \cdot \mathbf{k}}{k^2} \right\}. \quad (8) \end{aligned}$$

These relations can actually be obtained by taking derivatives with respect to the times η_j of the density consistency relations (1), using the second equality (7). As for the momentum consistency relations (4), these relations remain valid in the nonlinear regime and for biased small-scale fields $\tilde{\delta}_g(\mathbf{k}_j)$ and $\tilde{\lambda}_g(\mathbf{k}_j)$. The second term in Eq.(8), which arises from the $\tilde{\lambda}$ fields only, remains nonzero at equal times. This is due to the fact that λ involves the velocity or the time-derivative of the density, which probes the evolution between (infinitesimally close) different times.

The simplest relation associated with Eq.(8) is the bispectrum among two density-contrast fields and one momentum-divergence field,

$$\begin{aligned} \langle \tilde{\delta}(\mathbf{k}, \eta) \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\lambda}_g(\mathbf{k}_2, \eta_2) \rangle'_{k \rightarrow 0} &= -P_L(k, \eta) \frac{\mathbf{k}_1 \cdot \mathbf{k}}{k^2} \\ &\times \left(\langle \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\lambda}_g(\mathbf{k}_2, \eta_2) \rangle' \frac{D(\eta_1) - D(\eta_2)}{D(\eta)} + \right. \\ &\quad \left. + \langle \tilde{\delta}_g(\mathbf{k}_1, \eta_1) \tilde{\delta}_g(\mathbf{k}_2, \eta_2) \rangle' \frac{1}{D(\eta)} \frac{dD}{d\eta}(\eta_2) \right). \quad (9) \end{aligned}$$

At equal times, Eq.(9) reads as

$$\langle \tilde{\delta}(\mathbf{k}) \tilde{\delta}_g(\mathbf{k}_1) \tilde{\lambda}_g(\mathbf{k}_2) \rangle'_{k \rightarrow 0} = -\frac{\mathbf{k}_1 \cdot \mathbf{k}}{k^2} \frac{d \ln D}{d\eta} P_L(k) P_g(k_1). \quad (10)$$

3. Observable quantities

To test cosmological scenarios with the consistency relations of large-scale structures we need to relate them with observable quantities. We describe in this section the observational probes that we consider in this paper. We use the galaxy numbers counts or the weak lensing convergence to probe the density field. To apply the momentum consistency relations (6) and (10), we use the ISW effect to probe the momentum divergence λ (more precisely the time derivative of the gravitational potential and matter density) and the kSZ effect to probe the momentum \mathbf{p} .

3.1. Galaxy number density contrast δ_g

From galaxy surveys we can typically measure the galaxy density contrast within some redshift bin, smoothed with some finite-size window on the sky,

$$\delta_g^s(\boldsymbol{\theta}) = \int d\theta' W_\Theta(|\boldsymbol{\theta}' - \boldsymbol{\theta}|) \int d\eta I_g(\eta) \delta_g[r, r\boldsymbol{\theta}'; \eta], \quad (11)$$

where $W_\Theta(|\boldsymbol{\theta}' - \boldsymbol{\theta}|)$ is a 2D symmetric window function centered on the direction $\boldsymbol{\theta}$ on the sky, of characteristic angular radius Θ , $I_g(\eta)$ is the radial weight along the line of sight associated with a normalized galaxy selection function $n_g(z)$,

$$I_g(\eta) = \left| \frac{dz}{d\eta} \right| n_g(z), \quad (12)$$

$r = \eta_0 - \eta$ is the radial comoving coordinate along the line of sight, and η_0 is the conformal time today. Here and in the following we use the flat sky approximation, and $\boldsymbol{\theta}$ is the 2D vector that describes the direction on the sky of a given line of sight. The superscript ‘‘s’’ in δ_g^s denotes that we smooth the galaxy density contrast with the finite-size window W_Θ . Expanding in Fourier space the galaxy density contrast we can write

$$\begin{aligned} \delta_g^s(\boldsymbol{\theta}) &= \int d\theta' W_\Theta(|\boldsymbol{\theta}' - \boldsymbol{\theta}|) \int d\eta I_g(\eta) \\ &\quad \times \int d\mathbf{k} e^{i\mathbf{k}_\parallel r + i\mathbf{k}_\perp \cdot r\boldsymbol{\theta}'} \tilde{\delta}_g(\mathbf{k}, \eta) \end{aligned} \quad (13)$$

where k_\parallel and \mathbf{k}_\perp are respectively the parallel and the perpendicular components of the 3D wavenumber $\mathbf{k} = (k_\parallel, \mathbf{k}_\perp)$ (with respect to the reference direction $\boldsymbol{\theta} = 0$, and we work in the small-angle limit $\theta \ll 1$). Defining the 2D Fourier transform of the window W_Θ as

$$\tilde{W}_\Theta(|\boldsymbol{\ell}|) = \int d\boldsymbol{\theta} e^{-i\boldsymbol{\ell} \cdot \boldsymbol{\theta}} W_\Theta(|\boldsymbol{\theta}|), \quad (14)$$

we obtain

$$\delta_g^s(\boldsymbol{\theta}) = \int d\eta I_g(\eta) \int d\mathbf{k} \tilde{W}_\Theta(k_\perp r) e^{i\mathbf{k}_\parallel r + i\mathbf{k}_\perp \cdot r\boldsymbol{\theta}} \tilde{\delta}_g(\mathbf{k}, \eta). \quad (15)$$

3.2. Weak lensing convergence κ

From weak lensing surveys we can measure the weak lensing convergence, given in the Born approximation by

$$\kappa^s(\boldsymbol{\theta}) = \int d\theta' W_\Theta(|\boldsymbol{\theta}' - \boldsymbol{\theta}|) \int d\eta r g(r) \nabla^2 \frac{\Psi + \Phi}{2} [r, r\boldsymbol{\theta}'; \eta], \quad (16)$$

where Ψ and Φ are the Newtonian gauge gravitational potentials and the kernel $g(r)$ that defines the radial depth of the survey is

$$g(r) = \int_r^\infty dr_s \frac{dz_s}{dr_s} n_g(z_s) \frac{r_s - r}{r_s}, \quad (17)$$

where $n_g(z_s)$ is the redshift distribution of the source galaxies. Assuming no anisotropic stress, i.e. $\Phi = \Psi$, and using the Poisson equation,

$$\nabla^2 \Psi = 4\pi \mathcal{G}_N \bar{\rho}_0 \delta / a, \quad (18)$$

where \mathcal{G}_N is the Newton constant, $\bar{\rho}_0$ is the mean matter density of the Universe today, and a is the scale factor, we obtain

$$\kappa^s(\boldsymbol{\theta}) = \int d\eta I_k(\eta) \int d\mathbf{k} \tilde{W}_\Theta(k_\perp r) e^{i\mathbf{k}_\parallel r + i\mathbf{k}_\perp \cdot r\boldsymbol{\theta}} \tilde{\delta}(\mathbf{k}, \eta), \quad (19)$$

with

$$I_k(\eta) = 4\pi \mathcal{G}_N \bar{\rho}_0 \frac{rg(r)}{a}. \quad (20)$$

3.3. ISW secondary anisotropy Δ_{ISW}

From Eq.(7) λ can be obtained from the momentum divergence or from the time derivative of the density contrast. These quantities are not as directly measured from galaxy surveys as density contrasts. However, we can relate the time derivative of the density contrast to the ISW effect, which involves the time derivative of the gravitational potential. Indeed, the secondary cosmic microwave background temperature anisotropy due to the integrated Sachs-Wolfe effect along the direction $\boldsymbol{\theta}$ reads as (Garriga et al. 2004)

$$\begin{aligned} \Delta_{\text{ISW}}(\boldsymbol{\theta}) &= \int d\eta e^{-\tau(\eta)} \left(\frac{\partial \Psi}{\partial \eta} + \frac{\partial \Phi}{\partial \eta} \right) [r, r\boldsymbol{\theta}; \eta] \\ &= 2 \int d\eta e^{-\tau(\eta)} \frac{\partial \Psi}{\partial \eta} [r, r\boldsymbol{\theta}; \eta], \end{aligned} \quad (21)$$

where $\tau(\eta)$ is the optical depth, which takes into account the possibility of late reionization, and in the second line we assumed no anisotropic stress, i.e. $\Phi = \Psi$. We can relate Δ_{ISW} to λ through the Poisson equation (18), which reads in Fourier space as

$$-k^2 \tilde{\Psi} = 4\pi \mathcal{G}_N \bar{\rho}_0 \tilde{\delta} / a. \quad (22)$$

This gives

$$\frac{\partial \tilde{\Psi}}{\partial \eta} = \frac{4\pi \mathcal{G}_N \bar{\rho}_0}{k^2 a} (\tilde{\lambda} + \mathcal{H} \tilde{\delta}), \quad (23)$$

where $\mathcal{H} = d \ln a / d\eta$ is the conformal expansion rate. Integrating the ISW effect δ_{ISW} over some finite-size window on the sky, we obtain as in Eq.(15)

$$\begin{aligned} \Delta_{\text{ISW}}^s(\boldsymbol{\theta}) &= \int d\eta I_{\text{ISW}}(\eta) \int d\mathbf{k} \tilde{W}_\Theta(k_\perp r) e^{i\mathbf{k}_\parallel r + i\mathbf{k}_\perp \cdot r\boldsymbol{\theta}} \\ &\quad \times \frac{\tilde{\lambda} + \mathcal{H} \tilde{\delta}}{k^2}, \end{aligned} \quad (24)$$

with

$$I_{\text{ISW}}(\eta) = 8\pi \mathcal{G}_N \bar{\rho}_0 \frac{e^{-\tau}}{a}. \quad (25)$$

3.4. Kinematic SZ secondary anisotropy Δ_{kSZ}

Thomson scattering of CMB photons off moving free electrons in the hot galactic or cluster gas generates secondary anisotropies (Sunyaev & Zeldovich 1980; Gruzinov & Hu 1998; Knox et al.

1998). The temperature perturbation, $\Delta_{\text{kSZ}} = \delta T / T$, due to this kinematic Sunyaev Zeldovich (kSZ) effect, is

$$\Delta_{\text{kSZ}}(\boldsymbol{\theta}) = - \int d\mathbf{l} \cdot \mathbf{v}_e \sigma_T n_e e^{-\tau} = \int d\eta I_{\text{kSZ}}(\eta) \mathbf{n}(\boldsymbol{\theta}) \cdot \mathbf{p}_e, \quad (26)$$

where τ is again the optical depth, σ_T the Thomson cross section, \mathbf{l} the radial coordinate along the line of sight, n_e the number density of free electrons, \mathbf{v}_e their peculiar velocity, and $\mathbf{n}(\boldsymbol{\theta})$ the radial unit vector pointing to the line of sight. We also defined the kSZ kernel by

$$I_{\text{kSZ}}(\eta) = -\sigma_T \bar{n}_e a e^{-\tau}, \quad (27)$$

and the free electrons momentum \mathbf{p}_e as

$$n_e \mathbf{v}_e = \bar{n}_e (1 + \delta_e) \mathbf{v}_e = \bar{n}_e \mathbf{p}_e. \quad (28)$$

Because of the projection $\mathbf{n} \cdot \mathbf{p}_e$ along the line of sight, some care must be taken when we smooth $\Delta_{\text{kSZ}}(\boldsymbol{\theta})$ over some finite size angular window $W_\Theta(|\boldsymbol{\theta}' - \boldsymbol{\theta}|)$. Indeed, because the different lines of sight $\boldsymbol{\theta}'$ in the conical window are not perfectly parallel, if we define the longitudinal and transverse momentum components by the projection with respect to the mean line of sight $\mathbf{n}(\boldsymbol{\theta})$ of the circular window, e.g. $p_{e\parallel} = \mathbf{n}(\boldsymbol{\theta}) \cdot \mathbf{p}_e$, the projection $\mathbf{n}(\boldsymbol{\theta}') \cdot \mathbf{p}_e$ receives contributions from both $p_{e\parallel}$ and $\mathbf{p}_{e\perp}$. In the limit of small angles we could a priori neglect the contribution associated with $\mathbf{p}_{e\perp}$, which is multiplied by an angular factor and vanishes for a zero-size window. However, for small but finite angles, we need to keep this contribution because fluctuations along the lines of sight are damped by the radial integrations and vanish in the Limber approximation, which damps the contribution associated with $p_{e\parallel}$.

For small angles we write at linear order $\mathbf{n}(\boldsymbol{\theta}) = (\theta_x, \theta_y, 1)$, close to a reference direction $\boldsymbol{\theta} = 0$. Then, the integration over the angular window gives for the smoothed kSZ effect

$$\begin{aligned} \Delta_{\text{kSZ}}^s(\boldsymbol{\theta}) &= \int d\eta I_{\text{kSZ}}(\eta) \int d\mathbf{k} e^{i\mathbf{k} \cdot \mathbf{nr}} \left[\tilde{p}_{e\parallel} \tilde{W}_\Theta(k_\perp r) \right. \\ &\quad \left. - i \frac{\mathbf{k}_\perp \cdot \tilde{\mathbf{p}}_{e\perp}}{k_\perp} \tilde{W}'_\Theta(k_\perp r) \right]. \end{aligned} \quad (29)$$

Here we expressed the result in terms of the longitudinal and transverse components of the wave numbers and momenta with respect to the mean line of sight $\mathbf{n}(\boldsymbol{\theta})$ of the circular window W_Θ . Thus, whereas the radial unit vector is $\mathbf{n}(\boldsymbol{\theta}) = (\theta_x, \theta_y, 1)$, we can define the transverse unit vectors as $\mathbf{n}_{\perp x} = (1, 0, -\theta_x)$ and $\mathbf{n}_{\perp y} = (0, 1, -\theta_y)$, and we write for instance $\mathbf{k} = k_{\perp x} \mathbf{n}_{\perp x} + k_{\perp y} \mathbf{n}_{\perp y} + k_\parallel \mathbf{n}$. We denote $\tilde{W}'_\Theta(\ell) = d\tilde{W}_\Theta/d\ell$. The last term in Eq.(29) is due to the finite size Θ of the smoothing window, which makes the lines of sight within the conical beam not strictly parallel. It vanishes for an infinitesimal window, where $W_\Theta(\boldsymbol{\theta}) = \delta_D(\boldsymbol{\theta})$ and $\tilde{W}'_\Theta = 0$.

4. Consistency relation for the ISW temperature anisotropy

In this section we consider cross correlations with the ISW effect. This allows us to apply the consistency relation (9), which involves the momentum divergence λ and remains nonzero at equal times.

4.1. Galaxy-galaxy-ISW correlation

To take advantage of the consistency relation (9), we must consider three-point correlations ξ_3 (in configuration space) with one observable that involves the momentum divergence λ . Here, using the expression (24), we study the cross-correlation between two galaxy density contrasts and one ISW temperature anisotropy,

$$\xi_3(\delta_g^s, \delta_{g_1}^s, \Delta_{\text{ISW}_2}^s) = \langle \delta_g^s(\boldsymbol{\theta}) \delta_{g_1}^s(\boldsymbol{\theta}_1) \Delta_{\text{ISW}_2}^s(\boldsymbol{\theta}_2) \rangle. \quad (30)$$

The subscripts g , g_1 , and ISW_2 denote the three lines of sight associated with the three probes. Moreover, the subscripts g and g_1 recall that the two galaxy populations associated with δ_g^s and $\delta_{g_1}^s$ can be different and have different bias. As we recalled in section 2, the consistency relations rely on the undistorted motion of small-scale structures by large-scale modes. This corresponds to the squeezed limit $k \rightarrow 0$ in the Fourier-space equations (1) and (8), which writes more precisely as

$$k \ll k_L, \quad k \ll k_j, \quad (31)$$

where k_L is the wavenumber associated with the transition between the linear and nonlinear regimes. The first condition ensures that $\tilde{\delta}(k)$ is in the linear regime, while the second condition ensures the hierarchy between the large-scale mode and the small-scale modes. In configuration space, these conditions correspond to

$$\Theta \gg \Theta_L, \quad \Theta \gg \Theta_j, \quad |\boldsymbol{\theta} - \boldsymbol{\theta}_j| \gg |\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2|. \quad (32)$$

The first condition ensures that $\delta_g^s(\boldsymbol{\theta})$ is in the linear regime, whereas the next two conditions ensure the hierarchy of scales.

The expressions (15) and (24) give

$$\begin{aligned} \xi_3 = & \int d\eta d\eta_1 d\eta_2 I_g(\eta) I_{g_1}(\eta_1) I_{\text{ISW}_2}(\eta_2) \\ & \times \int d\mathbf{k} d\mathbf{k}_1 d\mathbf{k}_2 \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r_1) \tilde{W}_{\Theta_2}(k_{2\perp} r_2) \\ & \times e^{i(k_{\parallel} r + k_{1\parallel} r_1 + k_{2\parallel} r_2 + \mathbf{k}_\perp \cdot r \boldsymbol{\theta} + \mathbf{k}_{1\perp} \cdot r_1 \boldsymbol{\theta}_1 + \mathbf{k}_{2\perp} \cdot r_2 \boldsymbol{\theta}_2)} \\ & \times \langle \tilde{\delta}_g(\mathbf{k}, \eta) \tilde{\delta}_{g_1}(\mathbf{k}_1, \eta_1) \frac{\tilde{\lambda}(\mathbf{k}_2, \eta_2) + \mathcal{H}_2 \tilde{\delta}(\mathbf{k}_2, \eta_2)}{k_2^2} \rangle. \end{aligned} \quad (33)$$

The configuration-space conditions (32) ensure that we satisfy the Fourier-space conditions (31) and that we can apply the consistency relations (2) and (9). This gives

$$\begin{aligned} \xi_3 = & - \int d\eta d\eta_1 d\eta_2 b_g(\eta) I_g(\eta) I_{g_1}(\eta_1) I_{\text{ISW}_2}(\eta_2) \\ & \times \int d\mathbf{k} d\mathbf{k}_1 d\mathbf{k}_2 \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r_1) \tilde{W}_{\Theta_2}(k_{2\perp} r_2) \\ & \times e^{i(k_{\parallel} r + k_{1\parallel} r_1 + k_{2\parallel} r_2 + \mathbf{k}_\perp \cdot r \boldsymbol{\theta} + \mathbf{k}_{1\perp} \cdot r_1 \boldsymbol{\theta}_1 + \mathbf{k}_{2\perp} \cdot r_2 \boldsymbol{\theta}_2)} \\ & \times P_L(k, \eta) \frac{\mathbf{k}_1 \cdot \mathbf{k}}{k^2} \delta_D(\mathbf{k} + \mathbf{k}_1 + \mathbf{k}_2) \\ & \times \left(\langle \tilde{\delta}_{g_1} \frac{\tilde{\lambda}_2 + \mathcal{H}_2 \tilde{\delta}_2}{k_2^2} \rangle' \frac{D(\eta_1) - D(\eta_2)}{D(\eta)} \right. \\ & \left. + \langle \tilde{\delta}_{g_1} \frac{\tilde{\delta}_2}{k_2^2} \rangle' \frac{1}{D(\eta)} \frac{dD}{d\eta}(\eta_2) \right). \end{aligned} \quad (34)$$

Here we assumed that on large scales the galaxy bias is linear,

$$k \rightarrow 0: \quad \tilde{\delta}_g(\mathbf{k}) = b_g(\eta) \tilde{\delta}(\mathbf{k}) + \tilde{\epsilon}(\mathbf{k}), \quad (35)$$

where $\tilde{\epsilon}$ is a stochastic component that represents shot noise and the effect of small-scale (e.g., baryonic) physics on galaxy formation. From the decomposition (35), it is uncorrelated with the large-scale density field (Hamaus et al. 2010), $\langle \tilde{\delta}(\mathbf{k}) \tilde{\epsilon}(\mathbf{k}) \rangle = 0$. Then, in Eq.(34) we neglected the term $\langle \tilde{\epsilon} \tilde{\delta}_{g_1} (\tilde{\lambda}_2 + \mathcal{H}_2 \tilde{\delta}_2) \rangle$. Indeed, the small-scale local processes within the region $\boldsymbol{\theta}$ should be very weakly correlated with the density fields in the distant regions $\boldsymbol{\theta}_1$ and $\boldsymbol{\theta}_2$, which at leading order are only sensitive to the total mass within the large-scale region $\boldsymbol{\theta}$. Therefore, $\langle \tilde{\epsilon} \tilde{\delta}_{g_1} (\tilde{\lambda}_2 + \mathcal{H}_2 \tilde{\delta}_2) \rangle$ should exhibit a fast decay at low k , whereas the term in Eq.(34) associated with the consistency relation only decays as $P_L(k)/k \sim k^{n_s-1}$ with $n_s \simeq 0.96$. In Eq.(34), we also assumed that the galaxy bias b_g goes to a constant at large scales, which is usually the case, but we could take into account a scale dependence [by keeping the factor $b_g(k, \eta)$ in the integral over k].

The small-scale two-point correlations $\langle 1 \cdot 2 \rangle'$ are dominated by contributions at almost equal times, $\eta_1 \simeq \eta_2$, as different redshifts would correspond to points that are separated by several Hubble radii along the lines of sight and density correlations are negligible beyond Hubble scales. Therefore, ξ_3 is dominated by the second term that does not vanish at equal times. The integrals along the lines of sight suppress the contributions from longitudinal wavelengths below the Hubble radius c/H , while the angular windows only suppress the wavelengths below the transverse radii $c\Theta/H$. Then, for small angular windows, $\Theta \ll 1$, we can use Limber's approximation, $k_{\parallel} \ll k_{\perp}$ hence $k \simeq k_{\perp}$. Integrating over k_{\parallel} through the Dirac factor $\delta_D(k_{\parallel} + k_{1\parallel} + k_{2\parallel})$, and next over $k_{1\parallel}$ and $k_{2\parallel}$, we obtain the Dirac factors $(2\pi)^2 \delta_D(r_1 - r) \delta_D(r_2 - r)$. This allows us to integrate over η_1 and η_2 and we obtain

$$\begin{aligned} \xi_3 = & -(2\pi)^2 \int d\eta b_g(\eta) I_g(\eta) I_{g_1}(\eta) I_{\text{ISW}_2}(\eta) \frac{d \ln D}{d\eta} \\ & \times \int d\mathbf{k}_\perp d\mathbf{k}_{1\perp} d\mathbf{k}_{2\perp} \delta_D(\mathbf{k}_\perp + \mathbf{k}_{1\perp} + \mathbf{k}_{2\perp}) \tilde{W}_\Theta(k_\perp r) \\ & \times \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{2\perp} r) e^{i r(\mathbf{k}_\perp \cdot \boldsymbol{\theta} + \mathbf{k}_{1\perp} \cdot \boldsymbol{\theta}_1 + \mathbf{k}_{2\perp} \cdot \boldsymbol{\theta}_2)} \\ & \times P_L(k_\perp, \eta) \frac{\mathbf{k}_{1\perp} \cdot \mathbf{k}_\perp}{k_\perp^2 k_{2\perp}^2} P_{g_1, m}(k_{1\perp}, \eta), \end{aligned} \quad (36)$$

where $P_{g_1, m}$ is the galaxy-matter power spectrum. The integration over $\mathbf{k}_{2\perp}$ gives

$$\begin{aligned} \xi_3 = & -(2\pi)^2 \int d\eta b_g I_g I_{g_1} I_{\text{ISW}_2} \frac{d \ln D}{d\eta} \int d\mathbf{k}_\perp d\mathbf{k}_{1\perp} \tilde{W}_\Theta(k_\perp r) \\ & \times \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) P_L(k_\perp, \eta) P_{g_1, m}(k_{1\perp}, \eta) \\ & \times e^{i r[\mathbf{k}_\perp \cdot (\boldsymbol{\theta} - \boldsymbol{\theta}_2) + \mathbf{k}_{1\perp} \cdot (\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2)]} \frac{\mathbf{k}_{1\perp} \cdot \mathbf{k}_\perp}{k_{1\perp}^2 k_\perp^2}, \end{aligned} \quad (37)$$

and the integration over the angles of \mathbf{k}_\perp and $\mathbf{k}_{1\perp}$ gives

$$\begin{aligned} \xi_3 = & \frac{(\boldsymbol{\theta} - \boldsymbol{\theta}_2) \cdot (\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2)}{|\boldsymbol{\theta} - \boldsymbol{\theta}_2| |\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2|} (2\pi)^4 \int d\eta b_g I_g I_{g_1} I_{\text{ISW}_2} \frac{d \ln D}{d\eta} \\ & \times \int_0^\infty dk_\perp dk_{1\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) \\ & \times P_L(k_\perp, \eta) P_{g_1, m}(k_{1\perp}, \eta) J_1(k_\perp r |\boldsymbol{\theta} - \boldsymbol{\theta}_2|) \\ & \times J_1(k_{1\perp} r |\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2|), \end{aligned} \quad (38)$$

where J_1 is the first-order Bessel function of the first kind.

As the expression (38) arises from the kinematic consistency relations, it expresses the response of the small-scale two-point correlation $\langle \delta_{g_1}^s(\boldsymbol{\theta}_1) \Delta_{\text{ISW}_2}^s(\boldsymbol{\theta}_2) \rangle$ to a change of the initial condition associated with the large-scale mode $\delta_g^s(\boldsymbol{\theta})$. The kinematic effect

given at the leading order by Eq.(38) is due to the uniform motion of the small-scale structures by the large-scale mode. This explains why the result (38) vanishes in the two following cases

1. $(\theta - \theta_2) \perp (\theta_1 - \theta_2)$. There is a nonzero response of $\langle \delta_1 \lambda_2 \rangle$ if there is a linear dependence on $\delta(\theta)$ of $\langle \delta_1 \lambda_2 \rangle$, so that its first derivative is nonzero. A positive (negative) $\delta(\theta)$ leads to a uniform motion at θ_2 towards (away from) θ , along the direction $(\theta - \theta_2)$. From the point of view of θ_1 and θ_2 , there is a reflection symmetry with respect to the axis $(\theta_1 - \theta_2)$. For instance, if $\delta_1 > 0$ the density contrast at a position θ_3 typically decreases in the mean with the radius $|\theta_3 - \theta_1|$, and for $\Delta\theta_2 \perp (\theta_1 - \theta_2)$ the points $\theta_3^\pm = \theta_2 \pm \Delta\theta_2$ are at the same distance from θ_1 and have the same density contrast δ_3 in the mean, with typically $\delta_3 < \delta_2$ as $|\theta_3^\pm - \theta_1| > |\theta_2 - \theta_1|$. Therefore, the large-scale flow along $(\theta - \theta_2)$ leads to a positive $\lambda_2 = -\Delta\delta_2/\Delta\eta_2$ independently of whether the matter moves towards or away from θ (here we took a finite deviation $\Delta\theta_2$). This means that the dependence of $\langle \delta_1 \lambda_2 \rangle$ on $\delta(\theta)$ is quadratic [it does not depend on the sign of $\delta(\theta)$] and the first-order response function vanishes. Then, the leading-order contribution to ξ_3 vanishes. [For infinitesimal deviation $\Delta\theta_2$ we have $\lambda_2 = -\partial\delta_2/\partial\eta_2 = 0$; by this symmetry, in the mean δ_2 is an extremum of the density contrast along the orthogonal direction to $(\theta_1 - \theta_2)$.]
2. $\theta_1 = \theta_2$. This is a particular case of the previous configuration. Again, by symmetry from the viewpoint of δ_1 , the two points $\delta(\theta_2 + \Delta\theta_2)$ and $\delta(\theta_2 - \Delta\theta_2)$ are equivalent and the mean response associated with the kinematic effect vanishes.

This also explains why Eq.(38) changes sign with $(\theta_1 - \theta_2)$ and $(\theta - \theta_2)$. Let us consider for simplicity the case where the three points are aligned and $\delta(\theta) > 0$, so that the large-scale flow points towards θ . We also take $\delta_1 > 0$, so that in the mean the density is peaked at θ_1 and decreases outwards. Let us take θ_2 close to θ_1 , on the decreasing radial slope, and on the other side of θ_1 than θ . Then, the large-scale flow moves matter at θ_2 towards θ_1 , so that the density at θ_2 at a slightly later time comes from more outward regions (with respect to the peak at θ_1) with a lower density. This means that $\lambda_2 = -\partial\delta_2/\partial\eta_2$ is positive so that $\xi_3 > 0$. This agrees with Eq.(38), as $(\theta - \theta_2) \cdot (\theta_1 - \theta_2) > 0$ in this geometry, and we assume the integrals over wavenumbers are dominated by the peaks of $J_1 > 0$. If we flip θ_2 to the other side of θ_1 , we find on the contrary that the large-scale flow brings higher-density regions to θ_2 , so that we have the change of signs $\lambda_2 < 0$ and $\xi_3 < 0$. The same arguments explain the change of sign with $(\theta - \theta_2)$. In fact, it is the relative direction between $(\theta - \theta_2)$ and $(\theta_1 - \theta_2)$ that matters, measured by the scalar product $(\theta - \theta_2) \cdot (\theta_1 - \theta_2)$.

This geometrical dependence of the leading-order contribution to ξ_3 could provide a simple test of the consistency relation, without even computing the explicit expression in the right-hand side of Eq.(38).

4.2. Three-point correlation in terms of a two-point correlation

The three-point correlation ξ_3 in Eq.(38) cannot be written as a product of two-point correlations because there is only one integral along the line of sight that is left. However, if the linear power spectrum $P_L(k, z)$ is already known, we may write ξ_3 in terms of some two-point correlation ξ_2 . For instance, the small-scale cross-correlation between one galaxy density contrast and one weak lensing convergence,

$$\xi_2(\delta_{g_1}^s, \kappa_2^s) = \langle \delta_{g_1}^s(\theta_1) \kappa_2^s(\theta_2) \rangle \quad (39)$$

reads as

$$\xi_2 = (2\pi)^2 \int d\eta I_{g_1} I_{\kappa_2} \int_0^\infty dk_{1\perp} k_{1\perp} \tilde{F}_{\Theta_1}(k_{1\perp} r) \times \tilde{F}_{\Theta_2}(k_{1\perp} r) J_0(k_{1\perp} r |\theta_1 - \theta_2|) P_{g_1, m}(k_{1\perp}), \quad (40)$$

where we again used Limber's approximation. Here we denoted the angular smoothing windows by \tilde{F} to distinguish ξ_2 from ξ_3 . Then, we can write

$$\xi_3 = \frac{(\theta - \theta_2) \cdot (\theta_1 - \theta_2)}{|\theta - \theta_2| |\theta_1 - \theta_2|} \xi_2, \quad (41)$$

if the angular windows of the two-point correlation are chosen such that

$$\begin{aligned} \tilde{F}_{\Theta_1}(k_{1\perp} r) \tilde{F}_{\Theta_2}(k_{1\perp} r) &= (2\pi)^2 \frac{I_g I_{\text{ISW}_2}}{I_{\kappa_2}} b_g \frac{d \ln D}{d\eta} \\ &\times \left(\int_0^\infty dk_\perp \tilde{W}_\Theta(k_\perp r) J_1(k_\perp r |\theta - \theta_2|) P_L(k_\perp, \eta) \right) \\ &\times \frac{\tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) J_1(k_{1\perp} r |\theta_1 - \theta_2|)}{k_{1\perp} J_0(k_{1\perp} r |\theta_1 - \theta_2|)}. \end{aligned} \quad (42)$$

This implies that the angular windows \tilde{F}_{Θ_1} and \tilde{F}_{Θ_2} of the two-point correlation ξ_2 have an explicit redshift dependence.

In practice, the expression (42) may not be very convenient. Then, to use the consistency relation (38) it may be more practical to first measure the power spectra P_L and $P_{g_1, m}$ independently, at the redshifts needed for the integral along the line of sight (38), and next compare the measure of ξ_3 with the expression (38) computed with these power spectra.

4.3. Lensing-lensing-ISW correlation

From Eq.(38) we can directly obtain the lensing-lensing-ISW three-point correlation,

$$\xi_3(\kappa^s, \kappa_1^s, \Delta_{\text{ISW}_2}^s) = \langle \kappa^s(\theta) \kappa_1^s(\theta_1) \Delta_{\text{ISW}_2}^s(\theta_2) \rangle, \quad (43)$$

by replacing the galaxy kernels $b_g I_g$ and I_{g_1} by the lensing convergence kernels I_κ and I_{κ_1} ,

$$\begin{aligned} \xi_3 &= \frac{(\theta - \theta_2) \cdot (\theta_1 - \theta_2)}{|\theta - \theta_2| |\theta_1 - \theta_2|} (2\pi)^4 \int d\eta I_\kappa I_{\kappa_1} I_{\text{ISW}_2} \frac{d \ln D}{d\eta} \\ &\times \int_0^\infty dk_\perp dk_{1\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) \\ &\times P_L(k_\perp, \eta) P(k_{1\perp}, \eta) J_1(k_\perp r |\theta - \theta_2|) \\ &\times J_1(k_{1\perp} r |\theta_1 - \theta_2|). \end{aligned} \quad (44)$$

As compared with Eq.(38), the advantage of the cross-correlation with the weak lensing convergence κ is that Eq.(44) involves the matter power spectrum $P(k_{1\perp})$ instead of the more complicated galaxy-matter cross power spectrum $P_{g_1, m}(k_{1\perp})$.

5. Consistency relation for the kSZ effect

In this section we consider cross correlations with the kSZ effect. This allows us to apply the consistency relation (5), which involves the momentum \mathbf{p} and remains nonzero at equal times.

5.1. Galaxy-galaxy-kSZ correlation

In a fashion similar to the galaxy-galaxy-ISW correlation studied in section 4.1, we consider the three-point correlation between two galaxy density contrasts and one kSZ CMB anisotropy,

$$\xi_3(\delta_g^s, \delta_{g_1}^s, \Delta_{\text{kSZ}_2}^s) = \langle \delta_g^s(\boldsymbol{\theta}) \delta_{g_1}^s(\boldsymbol{\theta}_1) \Delta_{\text{kSZ}_2}^s(\boldsymbol{\theta}_2) \rangle, \quad (45)$$

in the squeezed limit given by the conditions (31) in Fourier space and (32) in configuration space. The expressions (15) and (29) give

$$\xi_3 = \xi_{3\parallel} + \xi_{3\perp} \quad (46)$$

with

$$\begin{aligned} \xi_{3\parallel} = & \int d\eta d\eta_1 d\eta_2 I_g(\eta) I_{g_1}(\eta_1) I_{\text{kSZ}_2}(\eta_2) \int d\mathbf{k} d\mathbf{k}_1 d\mathbf{k}_2 \\ & \times e^{i(\mathbf{k}\cdot\mathbf{n}r + \mathbf{k}_1\cdot\mathbf{n}_1 r_1 + \mathbf{k}_2\cdot\mathbf{n}_2 r_2)} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r_1) \\ & \times \tilde{W}_{\Theta_2}(k_{2\perp} r_2) \langle \tilde{\delta}_g(\mathbf{k}, \eta) \tilde{\delta}_{g_1}(\mathbf{k}_1, \eta_1) \tilde{p}_{e\parallel}^{(\mathbf{n}_2)}(\mathbf{k}_2, \eta_2) \rangle \end{aligned} \quad (47)$$

and

$$\begin{aligned} \xi_{3\perp} = & -i \int d\eta d\eta_1 d\eta_2 I_g(\eta) I_{g_1}(\eta_1) I_{\text{kSZ}_2}(\eta_2) \int d\mathbf{k} d\mathbf{k}_1 d\mathbf{k}_2 \\ & \times e^{i(\mathbf{k}\cdot\mathbf{n}r + \mathbf{k}_1\cdot\mathbf{n}_1 r_1 + \mathbf{k}_2\cdot\mathbf{n}_2 r_2)} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r_1) \\ & \times \tilde{W}'_{\Theta_2}(k_{2\perp} r_2) \langle \tilde{\delta}_g(\mathbf{k}, \eta) \tilde{\delta}_{g_1}(\mathbf{k}_1, \eta_1) \frac{\mathbf{k}_{2\perp}^{(\mathbf{n}_2)} \cdot \tilde{\mathbf{p}}_{e\perp}^{(\mathbf{n}_2)}}{k_{2\perp}^{(\mathbf{n}_2)}}(\mathbf{k}_2, \eta_2) \rangle, \end{aligned} \quad (48)$$

where we split the longitudinal and transverse contributions to Eq.(29). Here $\{\mathbf{n}, \mathbf{n}_1, \mathbf{n}_2\}$ are the radial unit vectors that point to the centers $\{\boldsymbol{\theta}, \boldsymbol{\theta}_1, \boldsymbol{\theta}_2\}$ of the three circular windows, and $\{(k_\parallel^{(\mathbf{n})}, \mathbf{k}_\perp^{(\mathbf{n})}), (k_{1\parallel}^{(\mathbf{n}_1)}, \mathbf{k}_{1\perp}^{(\mathbf{n}_1)}), (k_{2\parallel}^{(\mathbf{n}_2)}, \mathbf{k}_{2\perp}^{(\mathbf{n}_2)})\}$ are the longitudinal and transverse wave numbers with respect to the associated central lines of sight [e.g., $k_\parallel^{(\mathbf{n})} = \mathbf{n} \cdot \mathbf{k}$].

The computation of the transverse contribution (48) is similar to the computation of the ISW three-point correlation (34), using again Limber's approximation. At lowest order we obtain

$$\begin{aligned} \xi_{3\perp} = & \frac{(\boldsymbol{\theta} - \boldsymbol{\theta}_1) \cdot (\boldsymbol{\theta}_2 - \boldsymbol{\theta}_1)}{|\boldsymbol{\theta} - \boldsymbol{\theta}_1| |\boldsymbol{\theta}_2 - \boldsymbol{\theta}_1|} (2\pi)^4 \int d\eta b_g I_g I_{g_1} I_{\text{kSZ}_2} \frac{d \ln D}{d\eta} \\ & \times \int_0^\infty dk_\perp dk_{2\perp} k_{2\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{2\perp} r) \tilde{W}'_{\Theta_2}(k_{2\perp} r) \\ & \times P_L(k_\perp, \eta) P_{g_1, e}(k_{2\perp}, \eta) J_1(k_\perp r |\boldsymbol{\theta} - \boldsymbol{\theta}_1|) \\ & \times J_1(k_{2\perp} r |\boldsymbol{\theta}_2 - \boldsymbol{\theta}_1|), \end{aligned} \quad (49)$$

where $P_{g_1, e}$ is the galaxy-free electrons cross power spectrum.

The computation of the longitudinal contribution (47) requires slightly more care. Applying the consistency relation (5) gives

$$\begin{aligned} \xi_{3\parallel} = & - \int d\eta d\eta_1 d\eta_2 b_g(\eta) I_g(\eta) I_{g_1}(\eta_1) I_{\text{kSZ}_2}(\eta_2) \\ & \times \int d\mathbf{k} d\mathbf{k}_1 d\mathbf{k}_2 \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r_1) \tilde{W}_{\Theta_2}(k_{2\perp} r_2) \\ & \times e^{i(\mathbf{k}\cdot\mathbf{n}r + \mathbf{k}_1\cdot\mathbf{n}_1 r_1 + \mathbf{k}_2\cdot\mathbf{n}_2 r_2)} D(\eta) P_{L0}(k) \frac{dD}{d\eta}(\eta_2) \\ & \times i \frac{\mathbf{n}_2 \cdot \mathbf{k}}{k^2} \langle \tilde{\delta}_{g_1} \tilde{\delta}_{e_2} \rangle' \delta_D(\mathbf{k} + \mathbf{k}_1 + \mathbf{k}_2), \end{aligned} \quad (50)$$

where we only kept the contribution that does not vanish at equal times, as it dominates the integrals along the lines of sight, and

we used $P_L(k, \eta) = D(\eta)^2 P_{L0}(k)$. If we approximate the three lines of sight as parallel, we can write $\mathbf{n}_2 \cdot \mathbf{k} = k_\parallel$, where the longitudinal and transverse directions coincide for the three lines of sight. Then, Limber's approximation, which corresponds to the limit where the radial integrations have a constant weight on the infinite real axis, gives a Dirac term $\delta_D(k_\parallel)$ and $\xi_{3\parallel} = 0$ [more precisely, as we recalled above Eq.(36), the radial integration gives $k_\parallel \lesssim H/c$ while the angular window gives $k_\perp \lesssim H/(c\Theta)$ so that $k_\parallel \ll k_\perp$]. Taking into account the small angles between the different lines of sight, as for the derivation of Eq.(29), the integration over \mathbf{k}_2 through the Dirac factor gives at leading order in the angles

$$\begin{aligned} \xi_{3\parallel} = & - \int d\eta d\eta_1 d\eta_2 b_g(\eta) I_g(\eta) D(\eta) I_{g_1}(\eta_1) I_{\text{kSZ}_2}(\eta_2) \frac{dD}{d\eta}(\eta_2) \\ & \times \int dk_\parallel d\mathbf{k}_\perp dk_{1\parallel} d\mathbf{k}_{1\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r_1) \tilde{W}_{\Theta_2}(k_{1\perp} r_2) \\ & \times e^{i[k_\parallel(r-r_2) + \mathbf{k}_\perp \cdot (\boldsymbol{\theta} - \boldsymbol{\theta}_2) r_2 + k_{1\parallel}(r_1-r_2) + \mathbf{k}_{1\perp} \cdot (\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2) r_2]} \\ & \times P_{L0}(k_\perp) P_{g_1, e}(k_{1\perp}; \eta_1, \eta_2) i \frac{k_\parallel + \mathbf{k}_\perp \cdot (\boldsymbol{\theta}_2 - \boldsymbol{\theta})}{k_\perp^2}. \end{aligned} \quad (51)$$

We used Limber's approximation to write for instance $P_{L0}(k) \simeq P_{L0}(k_\perp)$, but we kept the factor k_\parallel in the last term, as the transverse factor $\mathbf{k}_\perp \cdot (\boldsymbol{\theta}_2 - \boldsymbol{\theta})$, due to the small angle between the lines of sight \mathbf{n} and \mathbf{n}_2 , is suppressed by the small angle $|\boldsymbol{\theta}_2 - \boldsymbol{\theta}|$. We again split $\xi_{3\parallel}$ over two contributions, $\xi_{3\parallel} = \xi_{3\parallel}^{\parallel} + \xi_{3\parallel}^{\perp}$, associated with the factors k_\parallel and $\mathbf{k}_\perp \cdot (\boldsymbol{\theta}_2 - \boldsymbol{\theta})$ of the last term. Let us first consider the contribution $\xi_{3\parallel}^{\parallel}$. Writing $i k_\parallel e^{i k_\parallel (r-r_2)} = \frac{\partial}{\partial r} e^{i k_\parallel (r-r_2)}$, we integrate by parts over η . For simplicity we assume that the galaxy selection function I_g vanishes at $z = 0$,

$$I_g(\eta_0) = 0, \quad (52)$$

so that the boundary term at $z = 0$ vanishes. Then, the integrations over k_\parallel and $k_{1\parallel}$ give a factor $(2\pi)^2 \delta_D(r-r_2) \delta_D(r_1-r_2)$, and we can integrate over η and η_1 . Finally, the integration over the angles of the transverse wave numbers yields

$$\begin{aligned} \xi_{3\parallel}^{\parallel} = & -(2\pi)^4 \int d\eta \frac{d}{d\eta} [b_g I_g D] I_{g_1} I_{\text{kSZ}_2} \frac{dD}{d\eta} \\ & \times \int_0^\infty dk_\perp dk_{1\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) \\ & \times \frac{k_{1\perp}}{k_\perp} P_{L0}(k_\perp) P_{g_1, e}(k_{1\perp}, \eta) J_0(k_\perp r |\boldsymbol{\theta} - \boldsymbol{\theta}_2|) \\ & \times J_0(k_{1\perp} r |\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2|), \end{aligned} \quad (53)$$

where J_0 is the zeroth-order Bessel function of the first kind. For the transverse contribution $\xi_{3\parallel}^{\perp}$ we can proceed in the same fashion, without integration by parts over η . This gives

$$\begin{aligned} \xi_{3\parallel}^{\perp} = & -(2\pi)^4 \int d\eta b_g I_g I_{g_1} I_{\text{kSZ}_2} D \frac{dD}{d\eta} \\ & \times \int_0^\infty dk_\perp dk_{1\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) \\ & \times k_{1\perp} P_{L0}(k_\perp) P_{g_1, e}(k_{1\perp}, \eta) |\boldsymbol{\theta} - \boldsymbol{\theta}_2| J_1(k_\perp r |\boldsymbol{\theta} - \boldsymbol{\theta}_2|) \\ & \times J_0(k_{1\perp} r |\boldsymbol{\theta}_1 - \boldsymbol{\theta}_2|). \end{aligned} \quad (54)$$

Comparing Eq.(54) with Eq.(53), we find $\xi_{3\parallel}^{\perp} / \xi_{3\parallel}^{\parallel} \sim k_\perp r |\boldsymbol{\theta} - \boldsymbol{\theta}_2|$. If the cutoff on k_\perp is set by the Bessel functions, we obtain $\xi_{3\parallel}^{\perp} \sim \xi_{3\parallel}^{\parallel}$. For very small angles, $|\boldsymbol{\theta} - \boldsymbol{\theta}_2| \rightarrow 0$, the cutoff over k is set by the angular window $\tilde{W}_\Theta(k_\perp r)$ or by the falloff of the linear power spectrum $P_{L0}(k_\perp)$, and $\xi_{3\parallel}^{\perp} \ll \xi_{3\parallel}^{\parallel}$.

In contrast with Eq.(38), the kSZ three-point correlation, given by the sum of Eqs.(49), (53) and (54), does not vanish for orthogonal directions between the small-scale separation $(\theta_1 - \theta_2)$ and the large-scale separation $(\theta - \theta_2)$. Indeed, the leading order contribution in the squeezed limit to the response of $\langle \delta_1 \mathbf{p}_2 \rangle$ to a large-scale perturbation δ factors out as $\langle \delta_1 \delta_2 \rangle \mathbf{v}_2$, where we only take into account the contribution that does not vanish at equal times (and we discard the finite-size smoothing effects). The intrinsic small-scale correlation $\langle \delta_1 \delta_2 \rangle$ does not depend on the large-scale mode δ , whereas \mathbf{v}_2 is the almost uniform velocity due to the large-scale mode, which only depends on the direction to $\delta(\theta)$ and is independent of the orientation of the small-scale mode $(\theta_1 - \theta_2)$.

Because the measurement of the kSZ effect only probes the radial velocity of the free electrons gas along the line of sight, which is generated by density fluctuations almost parallel to the line of sight over which we integrate and are damped by this radial integration, the result (53) is suppressed as compared with the ISW result (38) by the radial derivative $d \ln(b_g I_g D) / d\eta \sim 1/r$. Also, the contribution (53), associated with transverse fluctuations that are almost orthogonal to the second line of sight, is suppressed as compared with the ISW result (38) by the small angle $|\theta - \theta_2|$ between the two lines of sight.

One drawback of the kSZ consistency relation, (49) and (53)-(54), is that it is not easy to independently measure the galaxy-free electrons power spectrum $P_{g_1, e}$, which is needed if we wish to test this relation. Alternatively, Eqs.(53)-(54) may be used as a test of models for the free electrons distribution and the cross power spectrum $P_{g_1, e}$.

5.2. Lensing-lensing-kSZ correlation

Again, from Eqs.(49) and (53)-(54) we can directly obtain the lensing-lensing-kSZ three-point correlation,

$$\xi_3(\kappa^s, \kappa_1^s, \Delta_{\text{kSZ}_2}^s) = \langle \kappa^s(\theta) \kappa_1^s(\theta_1) \Delta_{\text{kSZ}_2}^s(\theta_2) \rangle, \quad (55)$$

by replacing the galaxy kernels $b_g I_g$ and I_{g_1} by the lensing convergence kernels I_κ and I_{κ_1} . This gives $\xi_3 = \xi_{3\perp} + \xi_{3\parallel}^{\parallel} + \xi_{3\parallel}^{\perp}$ with

$$\begin{aligned} \xi_{3\perp} &= \frac{(\theta - \theta_1) \cdot (\theta_2 - \theta_1)}{|\theta - \theta_1| |\theta_2 - \theta_1|} (2\pi)^4 \int d\eta I_\kappa I_{\kappa_1} I_{\text{kSZ}_2} \frac{d \ln D}{d\eta} \\ &\times \int_0^\infty dk_\perp dk_{2\perp} k_{2\perp} \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{2\perp} r) \tilde{W}'_{\Theta_2}(k_{2\perp} r) \\ &\times P_L(k_\perp, \eta) P_{m, e}(k_{2\perp}, \eta) J_1(k_\perp r |\theta - \theta_1|) \\ &\times J_1(k_{2\perp} r |\theta_2 - \theta_1|), \end{aligned} \quad (56)$$

$$\begin{aligned} \xi_{3\parallel}^{\parallel} &= -(2\pi)^4 \int d\eta \frac{d}{d\eta} [I_\kappa D] I_{\kappa_1} I_{\text{kSZ}_2} \frac{dD}{d\eta} \int_0^\infty dk_\perp dk_{1\perp} \\ &\times \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) \frac{k_{1\perp}}{k_\perp} P_{L0}(k_\perp) \\ &\times P_{m, e}(k_{1\perp}, \eta) J_0(k_\perp r |\theta - \theta_2|) J_0(k_{1\perp} r |\theta_1 - \theta_2|), \end{aligned} \quad (57)$$

and

$$\begin{aligned} \xi_{3\parallel}^{\perp} &= -(2\pi)^4 \int d\eta I_\kappa I_{\kappa_1} I_{\text{kSZ}_2} D \frac{dD}{d\eta} \int_0^\infty dk_\perp dk_{1\perp} \\ &\times \tilde{W}_\Theta(k_\perp r) \tilde{W}_{\Theta_1}(k_{1\perp} r) \tilde{W}_{\Theta_2}(k_{1\perp} r) k_{1\perp} P_{L0}(k_\perp) \\ &\times P_{m, e}(k_{1\perp}, \eta) |\theta - \theta_2| J_1(k_\perp r |\theta - \theta_2|) J_0(k_{1\perp} r |\theta_1 - \theta_2|). \end{aligned} \quad (58)$$

This now involves the matter-free electrons cross power spectrum $P_{m, e}$.

The application of the relations above is, unfortunately, a nontrivial task in terms of observations: to test those relations one would require the mixed galaxy (matter) - free electrons power spectrum. One possibility would be to do a stacking analysis of several X-rays observations of the hot ionised gas by measuring the bremsstrahlung effect. For instance, one could infer $n_e n_p T^{-1/2}$, by making some reasonable assumptions about the plasma state, as performed in Fraser-McKelvie et al. (2011), with the aim to measure n_e in filaments. We would of course need to cover a large range of scales. For kpc scales, inside galaxies and in the intergalactic medium, one could use for instance silicon emission line ratios (Kwitter & Henry 1998; Henry et al. 1996). For Mpc scales, or clusters, one may use the SZ effect (Rossetti et al. 2016). Nevertheless, all these proposed approaches are quite speculative at this stage.

6. Conclusions

In this paper, we have shown how to relate the large-scale consistency relations with observational probes. Assuming the standard cosmological model (more specifically, the equivalence principle and Gaussian initial conditions), nonzero equal-times consistency relations involve the cross-correlations between galaxy or matter density fields with the velocity, momentum or time-derivative density fields. We have shown that these relations can be related to actual measurements by considering the ISW and kSZ effects, which indeed involve the time derivative of the matter density field and the free electrons momentum field. We focused on the lowest-order relations, which apply to three-point correlation functions or bispectra, because higher-order correlations are increasingly difficult to measure.

The most practical relation obtained in this paper is probably the one associated with the ISW effect, more particularly its cross-correlation with two cosmic weak lensing convergence statistics. Indeed, it allows one to write this three-point correlation function in terms of two matter density field power spectra (linear and nonlinear), which can be directly measured (e.g., by two-point weak lensing statistics). Moreover, the result, which is the leading-order contribution in the squeezed limit, shows a specific angular dependence as a function of the relative angular positions of the three smoothed observed statistics. Then, both the angular dependence and the quantitative prediction provide a test of the consistency relation, that is, of the equivalence principle and of primordial Gaussianity. If we consider instead the cross-correlation of the ISW effect with two galaxy density fields, we obtain a similar relation but it now involves the mixed galaxy-matter density power spectrum $P_{g, m}$ and the large-scale galaxy bias b_g . These two quantities can again be measured (e.g., by two-point galaxy-weak lensing statistics) and provide another test of the consistency relation.

The relations obtained with the kSZ effect are more intricate. They do not show a simple angular dependence, which would provide a simple signature, and they involve the galaxy-free electrons or matter-free electrons power spectra. These power spectra are more difficult to measure. One can estimate the free electron density in specific regions, such as filaments or clusters, through X-ray or SZ observations, or around typical structures by stacking analysis of clusters. This could provide an estimate of the free electrons cross power spectra and a check of the consistency relations. Although we can expect significant error bars, it would be interesting to check that the results remain consistent with the theoretical predictions.

A violation of these consistency relations would signal either a modification of gravity on cosmological scales or non-

Gaussian initial conditions. We leave to future works the derivation of the deviations associated with various nonstandard scenarios.

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