Combined constraints on deviations of dark energy from an ideal fluid from *Euclid* and *Planck*

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Accepted 2015 November 7. Received 2015 October 26; in original form 2015 June 23

ABSTRACT

Cosmological fluids are commonly assumed to be distributed in a spatially homogeneous way, while their internal properties are described by a perfect fluid. As such, they influence the Hubble expansion through their respective densities and equation-of-state parameters. The subject of this paper is an investigation of the fluid-mechanical properties of a dark energy fluid, which is characterized by its sound speed and its viscosity apart from its equation of state. In particular, we compute the predicted spectra for the integrated Sachs–Wolfe effect for our generalized fluid, and compare them with the corresponding predictions for weak gravitational lensing and galaxy clustering, which had been computed in previous work. We perform statistical forecasts and show that the integrated Sachs–Wolfe signal obtained by cross-correlating *Euclid* galaxies with *Planck* temperatures, when joined to galaxy clustering and weak lensing observations, yields a percent sensitivity on the dark energy sound speed and viscosity. We prove that the iSW effect provides strong degeneracy breaking for low sound speeds and large differences between the sound speed and viscosity parameters.

Key words: gravitational lensing: weak — methods: analytical — cosmic background radiation — cosmology: theory — large-scale structure of Universe.

1 INTRODUCTION

The expansion dynamics of the Universe is usually described by assuming (i) general relativity as the theory of gravity, (ii) a high degree of symmetry, namely spatial isotropy and homogeneity at each instant in time, and (iii) ideal fluids which source the gravitational fields. These three assumptions lead to the Friedmann equations for the time-evolution of the scalefactor and on the growth of structure. In particular, we compute the predicted spectra for the integrated Sachs–Wolfe effect for our generalized fluid, and compare them with the corresponding predictions for weak gravitational lensing and galaxy clustering, which had been computed in previous work. We perform statistical forecasts and show that the integrated Sachs–Wolfe signal obtained by cross-correlating *Euclid* galaxies with *Planck* temperatures, when joined to galaxy clustering and weak lensing observations, yields a percent sensitivity on the dark energy sound speed and viscosity. We prove that the iSW effect provides strong degeneracy breaking for low sound speeds and large differences between the sound speed and viscosity parameters.

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1 INTRODUCTION

The expansion dynamics of the Universe is usually described by assuming (i) general relativity as the theory of gravity, (ii) a high degree of symmetry, namely spatial isotropy and homogeneity at each instant in time, and (iii) ideal fluids which source the gravitational fields. These three assumptions lead to the Friedmann equations for the time-evolution of the scalefactor $a(t)$, which reflect the fact that Einstein’s field equations are of second order, and show acceleration or deceleration as phenomena.

The inclusion of the cosmological constant on grounds of Lovelock’s theorem, which states that the GR field equations are the most general ones in four dimensions, that derivatives of the metric need to be included up to second order, that the energy–moment tensor is conserved, yields a natural way to explain cosmic acceleration at late times.

Alternatively, one can introduce dynamical dark energy components based on scalar self-interacting fields, and interpret the energy–momentum tensor with the corresponding conservation law. This allows the identification of the homogeneous and isotropic field with a relativistic ideal fluid, whose relation between pressure and density is parametrized by an equation of state $w = p/\rho$. At the same time, this equation of state is the only free function that is allowed by Einstein’s field equations with the symmetry assumptions of the Robertson–Walker metric.

For these reasons, a central goal of cosmology is to investigate dark energy and the cosmological constant through their influence on the dynamics of the scalefactor and on the growth of structures. The fluid picture is attractive due to its generality. Apart from actual substances like relativistic components ($w = +1/3$) and non-relativistic components ($w = 0$), it is general enough to describe spatial curvature ($w = -1/3$) and the cosmological constant ($w = -1$). Moreover, the isotropy and homogeneity of the fluid ensure the Friedmann symmetries.

Dark energy models based on self-interacting scalar fields show a natural variation of the dark energy equation-of-state parameter. This is because their time evolution is governed by the Klein–Gordon equation, and therefore the kinetic and potential terms in their energy–momentum tensor evolve, leading to a time evolution in the equation of state. Therefore, their influence on the expansion dynamics of the Universe will vary in time. In the slow-roll limit one recovers values of $w$ close to $-1$, resulting in accelerated expansion.

Adopting a fluid picture is an important test of whether the dark energy fluid is ideal or not. If the fluid has inhomogeneities: (i) pressure fluctuations and density fluctuations are related through a sound speed $c_s^2$, (which in the most straightforward case describes...
an adiabatic compression of the fluid), (ii) anisotropic stresses might arise, and finally (iii) velocity perturbations can experience viscous forces that dissipate kinetic energy. (For literature in this field we refer to Battye & Moss 2006, 2009; Mota et al. 2007; Calabrese et al. 2011; Ballesteros et al. 2012; Sapone & Majerotto 2012; Appleby, Linder & Weller 2013; Sapone et al. 2013; Dossett & Ishak 2013; Basse et al. 2014; Sawicki et al. 2013; Amendola et al. 2014; Chang & Xu 2014; Chang, Lu & Xu 2014; Cardona, Hollenstein & Kunz 2014; Pearson 2014; Ballesteros 2015). In addition, there can be a non-linear relation between pressure and density of a fluid, one example of which would be Chaplygin cosmologies (Bento, Bertolami & Sen 2002; Li & Xu 2014). Alternatively, a similar phenomenology could in principle be due to modifications in gravity rather than due to non-ideal fluids under general relativity (Kunz & Sapone 2007; Bertschinger & Zukin 2008; Silvestri 2009; Pogosian et al. 2010; Song et al. 2010; Leon & Saridakis 2011; Salas & Kunz 2011; Baker, Ferreira & Skordis 2013; Boubekeur et al. 2014).

These modifications break homogeneity on small scales. They require corresponding fluid equations for their time evolution, as well couplings to local gravitational fields. The latter enable interaction between the dark energy fluid, the dark matter and the baryonic component. Commonly, one observes a difference between the two metric potentials in the case of non-zero sound speeds and equations of state unequal to $-1$. Such difference can be probed by photons, and is relative to the motion of non-relativistic objects such as galaxies, which are only sensitive to a single metric potential.

In this paper, we investigate cosmological perturbations with a non-ideal dark energy fluid and aim to forecast constraints on its speed of sound $c_s$ and its viscosity from Euclid\footnote{http://www.euclid-ec.org/} (Laureijs et al. 2011) and Planck (Planck Collaboration XVI 2014b; Planck Collaboration XIII 2015). Specifically, we consider tomographic weak gravitational lensing (Ayaita, Schäfer & Weber 2012), galaxy clustering (DeDeo, Caldwell & Steinhardt 2003; Ballesteros & Lesgourgues 2010; de Putter, Huterer & Linder 2010; Archidiacono, Lopez-Honorez & Mena 2014). At a first sight, it would appear that choosing a dark energy equation of state too far from the cosmological constant value is incompatible with present constraints (Planck Collaboration XIII 2015). However, when including extra parameters such as the speed of sound and viscosities in the fluid, constraints become much more loose (See e.g. Mota et al. 2007; Archidiacono et al. 2014).

This article is structured as follows. We develop the necessary perturbation equations for non-ideal dark energy fluids and a suitable parametrization in Section 2 and discuss cosmological probes in Section 3, before computing forecasts on non-ideal dark energy properties in Section 4. We summarize our results in Section 5. As regards the forecasts, let us remark that some of the authors are currently involved in a code comparison effort taking place within the Euclid Collaboration Forecast project. We want to stress then that some of the numbers in this paper might or might not change once agreement between different codes will be reached. The reference cosmological model is a spatially flat, dark-energy dominated model with the parameter choices \{\Omega_{m0}h^2, \Omega_b h^2, n_s, \Omega_m, w\} = \{0.142, 0.022, 0.96, 0.32, -0.8\}. The first four correspond to the constraints from Planck (Planck Collaboration XVI 2014b) and WMAP polarization low-multipole likelihood (Bennett et al. 2013; Planck Collaboration XV 2014a), and they represent the present official baseline for Euclid forecasts. The dark energy equation-of-state parameter was instead set to a fiducial value of $w = -0.8$, in order for it to be sufficiently far from $-1$ (see discussion above). The amplitude of the primordial power spectrum was fixed to $A_s = 2.1 \times 10^{-9}$.

## 2 COSMOLOGY WITH NON-IDEAL FLUIDS

### 2.1 Expansion dynamics

Since we focus on late cosmological times, where dark matter and dark energy are dominating the energy density of the Universe, we can approximate the Hubble function \( H(a) = \dot{a}/a \) with

\[
H^2 = H_0^2 \left[ \Omega_{m,0} a^{-3} + (1 - \Omega_{m,0}) a^{-3(1+w)} \right].
\]

(1)

Here, \( a \) is the scalefactor, \( \Omega_{m,0} \) is the dark matter density parameter today, \( w \) is the equation of state of dark energy, which we assume to be constant, and \( H_0 \) is the Hubble parameter today. In addition, we do not consider global curvature. The comoving distance is defined as

\[
\chi = c \int_a^1 \frac{da}{a^2 H(a)} = \chi_H \int_a^1 \frac{da}{a^2 H(a)/H_0},
\]

(2)

with the Hubble distance \( \chi_H = c/H_0 \simeq 2996.9 \text{ Mpc}/h \). At the same time, this defines conformal time \( \tau \) through \( \chi = c \tau \). In the following, we will set \( c = 1 \).

### 2.2 Perturbations and their analytical solutions

If we consider a non-ideal fluid dark energy, characterized by a constant equation of state \( w \), a speed of sound \( c_s \), and an anisotropic stress component \( \sigma \), we can write the evolution of \( \sigma \) as in Hu (1998):

\[
\sigma' + \frac{3}{a} \sigma = \frac{8}{3} c_s^2 \frac{V}{(1+w)^2 a^3 H^2}.
\]

(3)

Here, the prime indicates derivative with respect to \( a \) and \( c_s^2 \) is called viscosity parameter, as it gives a measure of the fluid’s viscosity. Indeed, equation (3) implies that when \( c_s^2 = 0 \), then the anisotropic
stress component $\sigma$ is also vanishing, while when e.g. $c_i^2 = 1/3$ the evolution of anisotropic stress for radiation up to the quadrupole is recovered.

To this equation, we add the first order perturbation equations for the density contrast $\delta$ and the velocity perturbation $V$

$$\delta' = 3(1 + w)\phi' - \frac{V}{Ha^2} - \frac{1}{3a} \left( \frac{\delta p}{\rho} - w\delta \right), \quad (4)$$

$$V' = -(1 - 3w)\frac{V}{a} + \frac{k^2}{Ha^2} \frac{\delta p}{\rho} + (1 + w)\frac{k^2}{Ha^2} \psi' + \frac{k^2}{Ha^2} \sigma, \quad (5)$$

where $\delta p$ is the pressure perturbation, $\rho$ is the dark energy density, $\psi$ and $\phi$ are the metric perturbations in the Newtonian gauge, defined by the line element

$$dx^2 = a^2 \left[ -(1 + 2\psi) d\tau^2 + (1 - 2\phi) dx_i dx^i \right]. \quad (6)$$

Pressure perturbations are parametrized as

$$\delta p = c_i^2 \rho s + \frac{3aH(c_i^2 - c_s^2)}{k^2} \rho V, \quad (7)$$

where $c_i^2 \equiv \rho / \rho = w$ is the adiabatic speed of sound for a fluid with constant equation of state. $c_i^2$ reduces to $c_s^2$ in the case of a perfect fluid, when no dissipative effects are present, which would otherwise lead to entropic perturbations (Bean & Doré 2004).

In order to close the differential equation system, one needs to include the Poisson equation

$$k^2\phi = -4\pi Ga^2 \sum_i \rho_i \left( \delta_i + \frac{3aH}{k^2} V_i \right) = -4\pi Ga^2 \sum_i \rho_i \Delta_i, \quad (8)$$

where the sum runs over all clustering fluids, $G$ is the Newton constant, and in the last equality we have defined the gauge-invariant density perturbation of the $i$th fluid, $\Delta_i \equiv \delta_i + 3aHV_i/k^2$ and the fourth Einstein equation

$$k^2 (\phi - \psi) = 12\pi Ga^2 (1 + w) \rho \sigma \quad (9)$$

$$= \frac{9}{2} H_0^2 (1 - \Omega_m,0) a^{-1(3+w)} (1 + w) \sigma$$

$$\equiv B(a) \sigma. \quad (10)$$

In Sapone & Majerotto (2012) the following analytical solutions for $\delta$, $V$ and $\sigma$ were found for the matter dominated era:

$$\delta = \frac{3(1 + w)^2}{3c_i^2(1 + w) + 8(c_i^2 - w)c_i^2 k^2}, \quad (11)$$

$$V = -\frac{9(1 + w)^2 (c_i^2 - w)}{3c_i^2(1 + w) + 8c_i^2(c_i^2 - w)} H_0 \sqrt{\Omega_m,0} \frac{\phi_0}{\sqrt{ak^2}},$$

$$= -3aH \frac{c_i^2}{c_i^2 - w} \delta, \quad (12)$$

$$\sigma = -\frac{8c_i^2 (c_i^2 - w)}{3c_i^2(1 + w) + 8(c_i^2 - w)c_i^2 k^2}, \quad (13)$$

where $k^2 \phi \approx -\phi_0$. The latter is valid strictly only during matter domination and while neglecting dark energy perturbations. Nevertheless, the validity of this approximation extends to redshifts very close to today’s $z$, as can be seen from fig. 1 of Sapone & Majerotto (2012).

As found in Sapone & Majerotto (2012), to which we refer for further detail on the analytic solutions, the relevant quantity is the effective sound speed

$$c_{\text{eff}}^2 = c_i^2 + \frac{8}{3} c_i^2 w - \frac{w}{1 + w}, \quad (14)$$

as equations (11)–(13) can be rewritten in terms of it. This means that the sound speed and the viscosity have a similar damping effect on density and velocity perturbations (as also noticed in Mota et al. 2007; Calabrese et al. 2011). It is interesting to notice that the effect of $c_i^2$ is enhanced with respect to that of $c_s^2$ by a factor of $8(c_i^2 - w)/[3(1 + w)]$, which is $\sim 10$ if $w \sim -0.8$ and $c_s^2$ is very small. A relative large $w$ and very small $c_s^2$ are precisely the cases where a viscosity can be observed best, as will be shown in the following sections. Moreover $c_{\text{eff}}^2$ is bounded by the case of a cosmological constant, as equation (11) diverges for $w = -1$.

### 2.3 Observable parameters

To understand how the viscosity affects the physical observables, it is useful to introduce the clustering parameter $Q$ and the anisotropy parameter $\eta$. These were defined in Amendola, Kunz & Sapone (2008) and computed in the case of viscous dark energy in Sapone & Majerotto (2012).

$Q$ parametrizes the deviation from a purely matter-dominated Newtonian potential and is given by (see Sapone & Majerotto 2012)

$$Q - 1 \equiv \frac{\rho \Delta}{\rho_m \Delta_m} = 1 - \frac{\Omega_m,0}{\Omega_m,0} (1 + w) \frac{a^{-3w}}{1 - 3w + \frac{2w}{3H_0^2 \Omega_m,0 c_{\text{eff}}^2}}$$

$$= \frac{Q_0}{1 + \alpha a}, \quad (15)$$

where $\alpha = 2k^2 c_{\text{eff}}^2 / [(3H_0^2 \Omega_m,0)(1 - 3w)]$ and $Q_0 = (1 + w)(1 - \Omega_m,0)/(\Omega_m,0(1 - 3w))$.

The anisotropy parameter is then given by

$$\eta \equiv \frac{\psi}{\phi} - 1 = -\frac{9}{2} H_0^2 (1 - \Omega_m,0) (1 + w) a^{-1 - 3w} k^2 Q \left( 1 - \frac{c_i^2}{c_{\text{eff}}^2} \right). \quad (16)$$

This is non-zero only when anisotropic stress is present and the metric perturbations $\psi$ and $\phi$ are different.

Let us finally define the parameter $\Sigma$ (Amendola et al. 2008) as

$$\Sigma = \left( 1 + \frac{1}{2\eta} \right) Q. \quad (17)$$

This is useful because it represents the deviation of the weak lensing potential $\Phi = \psi + \phi$ from its behaviour in the case of no dark energy perturbations. It is important to remark that for all our forecasts we used the numerical solutions to equations (3)–(5) and (8)–(9) computed with a modified version of the {	extsc camb} code2 (Lewis, Challinor & Lasenby 2000), and not their analytic approximations, equations (11)–(13). We did use instead the analytical expression for the $Q$ and $\Sigma$ parameters. These approximate very well their numerical counterparts up to redshifts very close to 0, as shown in figs 3 and 4 of Sapone & Majerotto (2012).

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1. http://camb.info
3 COSMOLOGICAL PROBES

In Sapone et al. (2013), we forecasted constraints to the viscosity parameter and the sound speed from the Euclid galaxy clustering and weak lensing surveys. Here, we aim to complete the picture by adding to the latter the constraints from the iSW tomography signal obtained by cross-correlating galaxies mapped by the Euclid photometric instrument with the Planck temperature map. This provides a combination of all major probes of cosmic structure formation, which draw their sensitivity from the growth rate and interactions between fluids, from the shape of the initial perturbations and from the expansion history.

The Euclid survey is a mission of the ESA Cosmic Vision programme that will be launched in 2020, and will perform both a photometric and a spectroscopic survey, the first aiming mainly at measuring weak lensing while the second at measuring the galaxy power spectrum. The Planck satellite is also a mission of ESA Cosmic Vision programme, which mapped the CMB fluctuations with unprecedented precision and control of systematic effects.

To perform our forecasts, we use the Fisher matrix (Tegmark et al. 1997), which quantifies the decrease in likelihood if a model parameter moves away from the fiducial value, and can be computed for a local Gaussian approximation to the likelihood $L \propto \exp(-\chi^2/2)$. In our forecasts we assume the official Euclid specifications, that can be found in Laureijs et al. (2011). The fiducial cosmological parameters correspond to the 2013 Planck measurements (Planck Collaboration XVI 2014b), except for the value of $w$, for which we assume $w = -0.8$, in order for the effects to be more clearly visible (as done in Sapone & Majerotto 2012; Sapone et al. 2013) and of course except for the values of $c_s^2$ and $\epsilon_s^2$.

In the following, we will describe the iSW tomography signal and give a short summary on the signal coming from weak lensing and the galaxy power spectrum from spectroscopy.

3.1 iSW signal

When a CMB photon moves into a time-evolving metric such as that of equation (6) the unbalance between the blueshift experienced at the entrance and the redshift experienced at the exit of its varying potential well originates a perturbation $\zeta$ in the CMB temperature $T_{CMB}$ given by (Sachs & Wolfe 1967),

$$\zeta = \Delta T / T_{CMB} = \int d\tau \left( \frac{\partial \Phi}{\partial \tau} + \frac{\partial \psi}{\partial \tau} \right) = \int_0^{\tau_H} d\chi a^2 H a^2 \frac{\partial \Phi}{\partial a}, \quad (18)$$

where $\chi$ is the comoving distance (see equation 2) and $\Phi$ is the weak lensing potential.

In the case of pure matter domination, $\Phi = \text{const}$, hence the iSW effect vanishes, while in presence of any fluid with $w \neq 0$ the temperature fluctuation $\zeta$ will be non-zero. This is why the late iSW is particularly interesting to us, as it is originated by the appearing of dark energy and it is an independent proof of its existence (first detected by Boughn & Crittenden 2004; Giannantonio et al. 2012, for an updated measurement).

Let us now compute the term inside the integral, passing to Fourier space, and in the case of viscous dark energy (see also Schaefer 2009; Sapone & Majerotto 2012):

$$\frac{\partial \Phi}{\partial a} = -\frac{3}{2} \frac{H^2 \Omega_{m,0}}{a k^2} \left\{ \frac{\Sigma (a, k) \Delta_m (a, k)}{a} + \frac{\Sigma' (a, k) \Delta_m (a, k) - \frac{1}{a} \Sigma (a, k) \Delta_m (a, k)}{a} \right\}. \quad (19)$$

It is possible to see from this expression that anisotropic perturbations enter the iSW effect in two ways: by modifying $\Delta_m$ and through the additional presence of $\Sigma$ and $\Sigma'$. At linear order, it is possible to isolate today’s $\Delta_m$ from its time evolution:

$$\Delta_m (a, k) = a G (a, k) \Delta_m (a, k) \Delta_m (a, k), \quad (20)$$

where $\Delta_m (a, k) = \Delta_m (a = 1, k)$. We write hence equation (19) as:

$$\frac{\partial \Phi}{\partial a} = -\frac{3}{2} \frac{H^2 \Omega_{m,0}}{a k^2} \left( G (a, k) \Sigma (a, k) \Delta_m (a, k) \right), \quad (21)$$

and equation (18) reads now

$$\zeta = \int_0^{\tau_H} d\chi W_\chi (\chi \Delta_m (a, k)) \quad (22)$$

where the weighting function $W_\chi (\chi)$ is

$$W_\chi (\chi) = \frac{3}{2} \frac{H^2 \Omega_{m,0}}{a k^2} a^2 H \frac{\partial}{\partial a} G (a, k) \Sigma (a, k). \quad (23)$$

Since the iSW is a secondary effect of the CMB (Rees & Sciama 1968), it can be extracted through cross-correlation of the CMB temperature fluctuations to the galaxy density (Crittenden & Turok 1996). Let us hence write the galaxy density obtained through imaging surveys, in order to compute its cross-correlation with the iSW. The line-of-sight projected galaxy density $\gamma$ is given by (Smail et al. 1995)

$$\gamma = \int_0^{\tau_H} d\chi D(z) \frac{d\chi}{d\chi} b(\chi) G (\chi) k (z) \quad (24)$$

being $D(z)$ the galaxy distribution defined as

$$D(z) = \left( \frac{z}{z_0} \right)^2 \exp \left[ -\left( \frac{z}{z_0} \right)^{\beta_D} \right], \quad (25)$$

where the fiducial parameters $\beta_D$ and $z_0$ depend on the imaging survey considered. In the case of Euclid, they are $\beta_D = 3/2$, $z_0 = z_{\text{mean}}/\sqrt{2}$, and $z_{\text{mean}} = 0.9$ (Laureijs et al. 2011).

Even though the signal from the iSW increases noticeably when cross-correlating it with the galaxy density field, both the cross-correlation spectrum and the galaxy spectrum are line-of-sight integrated quantities, hence much information may be lost. For this reason, we decide to use iSW tomography (Douspis et al. 2008; Ho et al. 2008; Juergens & Schaefer 2012). In particular we divide the whole galaxy sample into 5 bins with equal number of galaxies (in order to match with the binning used by official Euclid documents for weak lensing tomography). To do this, we replace the galaxy distribution function $D(z)$, equation (25) in $\gamma$, equation (24), with the radial distribution function of galaxies in the $i$th bin $D_i(z)$.

The latter is obtained by binning the overall distribution $D(z)$ and convolving it with the photometric redshift distribution function ($\text{Amendola et al. 2008}$).

We are finally able to write our observable, i.e. the iSW-galaxy cross-correlation spectrum in the $i$-th redshift bin, along with the iSW-auto correlation spectrum $C_{\chi,\chi}(\ell)$ and the galaxy–galaxy auto correlation spectrum $C_{\gamma,\gamma,0}(\ell)$ of the $i$-jets, by applying a Limber projection (Limber 1954) in the flat-sky approximation:

$$C_{\chi,\chi}(\ell) = \int_0^{\tau_H} d\chi \frac{d\chi}{d\chi} W_\chi (\chi) W_\chi (\chi) P_{\Delta\Delta} (k = \ell/\chi), \quad (26)$$

$$C_{\chi,\chi}(\ell) = \int_0^{\tau_H} d\chi \frac{d\chi}{d\chi} W_\chi (\chi) P_{\Delta\Delta} (k = \ell/\chi) \quad (27)$$

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Here \( \tilde{P}_{\Delta i} (k) \) is the linear matter power spectrum today, and the galaxy weighting function of the \( i \)th bin \( W_{\gamma,i}(\chi) \) is

\[
W_{\gamma,i}(\chi) = D_i(z) \frac{dz}{d\chi} b(\chi) G(\chi),
\]

(29)

\( C_{\chi \gamma}(\ell) \) and \( C_{\chi \gamma 0}(\ell) \) have been computed because they are needed in order to estimate statistical errors of \( C_{\gamma \gamma}(\ell) \). The tomographic iSW spectra \( C_{\gamma \gamma}(\ell) \) are shown in Fig. 1 for two fiducial models: a standard dark energy model with \( c_s^2 = 1 \) and \( c_s^2 = 0 \) (solid lines) and a model with viscosity: \( c_s^2 = 10^{-3} \) and \( c_s^2 = 10^{-6} \) (dashed lines). The colour shading indicates the redshift bin for which \( C_{\gamma \gamma}(\ell) \) was evaluated. The iSW effect is a large-scale effect originating from low redshift, as the influence of dark energy on the growth of gravitational potentials in the large-scale structure is strongest. The effect of dark energy viscosity and small sound speed is strongest on large scales as well (see also Sapone & Majerotto 2012), and affects a wide range of multipoles. Keeping all cosmological parameters fixed, dark energy viscosity would increase the amplitude of the iSW effect by up to 25 per cent on large angular scales and at low redshift. This sensitivity of the spectra at low multipoles is fortunate because these scales can be well probed with the iSW effect.

### 3.2 Weak lensing

To the iSW tomography signal we add the weak lensing tomographic signal (Hu 1999, 2002; Heavens 2003; Jain & Taylor 2003), coming from the same photometric survey as \( \gamma \) and using the same redshift bins. Here, we only give the main equation expressing the weak lensing power spectrum, which is used for our forecasts, and refer to Sapone et al. (2013) for further details.

In presence of anisotropic stress, the weak lensing convergence power spectrum is given by (Hu 1999, 2002; Jain & Taylor 2003; Hu & Jain 2004)

\[
C_{\kappa \gamma ij}(\ell) = \int_0^{2\pi/\ell} \frac{d\chi}{\chi^2} W_{\gamma,i}(\chi) W_{\kappa,j}(\chi) \right|_{N}^\Sigma P_{\text{NL}}(k = \ell/\chi, \chi).
\]

(30)

where the subscript \( ij \) refers to the redshift bins around \( z_i \) and \( z_j \), with

\[
W_{\kappa,i}(\chi) = \frac{3\Omega_{m} F_i(\chi)}{2\chi H} \frac{1}{a} \chi
\]

(31)

\[
F_i(\chi) = \int_{0}^{\chi R} d\chi' n(\chi') D_i(\chi') \frac{\chi'}{\chi} - \chi
\]

(32)

and where \( D_i \) is the same tomographic distribution function of galaxies used for the iSW effect. While tomography in general greatly reduces statistical errors, the actual shape of the choice of the binning does not affect results in a serious way, although in principle there is room for optimization (Schäfer & Heisenberg 2012).

In Fig. 2 we show the tomographic weak lensing spectra \( C_{\kappa \gamma i}(\ell) \) for the same models and the same redshifts as in Fig. 1. As for iSW, the effect of viscosity is detected at large scales and for a large range of scales (but smaller than for iSW). Instead, contrary to iSW, here the sensitivity to viscosity is stronger at higher redshift. This is because the efficiency of weak lensing is higher for longer light paths.

In principle it is also possible to define a cross-spectrum of weak lensing and iSW, \( C_{\kappa \gamma i}(\ell) \), and of weak lensing and galaxy distribution, \( C_{\kappa \kappa i}(\ell) \), but both these spectra are subdominant with respect to \( C_{\gamma \gamma}(\ell) \). This is because the weak lensing convergence signal comes from the distortion of the light path at redshifts intermediate between us and the galaxies mapped by the imaging survey, while the iSW signal originates precisely at the same redshifts where the galaxies are. We have tested this fact by computing the signal to noise-ratio for measuring \( C_{\kappa \gamma i}(\ell) \) and found it much smaller than that of \( C_{\gamma \gamma}(\ell) \).

### 3.3 Spectroscopic galaxy power spectrum

To the iSW and weak lensing measurements, both measured through photometric observations, we add data coming from the power spectrum of spectroscopically observed galaxies. Here, we only show the expression of the observed power spectrum, which is needed in...
order to compute our forecasts, and refer again the reader to Sapone et al. (2013) for further detail.

Following Seo & Eisenstein (2003) we write the observed galaxy power spectrum as

\[ P_{\gamma\gamma}^\text{spec}(z, k, \mu) = D^2(z) H(z) G^2(z, k) b(z)^2 \left(1 + \beta \mu^2\right)^2 P_0(k) + \mathcal{P}_\text{bns}, \]

(33)

where the subscript r refers to the reference (or fiducial) cosmological model.

Here \( \mathcal{P}_\text{bns} \) is a scale-independent offset due to imperfect removal of shot-noise, \( \mu = k \cdot \hat{r} / k \), is the cosine of the angle of the wave mode with respect to the line of sight pointing into the direction \( \hat{r} \), \( P_0 \) is the fiducial matter power spectrum evaluated at redshift zero, \( G(z, k) \) is the linear growth factor of the matter perturbations, \( b(z) \) is the bias factor and \( D_A(z) \) is the angular diameter distance. The wavenumber \( k \) and \( \mu \) have also to be written in terms of the fiducial cosmology (see Seo & Eisenstein 2003; Amendola, Quercellini & Giallongo 2005; Sapone & Amendola 2007, for more details). The fiducial bias used can be found in Orsi et al. (2010), who derived their results by using a semi-analytical model of galaxy formation. The matter power spectrum has been computed with a modified version of the \textsc{CAMB} code accounting for anisotropies.

4 STATISTICAL ERRORS FORECASTS

In this section, we estimate marginalized statistical errors on the sound speed and viscosity parameters

\[ \frac{\partial \log \mathcal{C}_{\beta\gamma,i}}{\partial \theta_a}, \frac{\partial \log \mathcal{C}_{\beta\gamma,i}}{\partial \theta_b}. \]

(34)

Here the sum runs from \( \ell = 5 \) to \( \ell \in 300, \) \( \theta_a \) are the cosmological parameters, \( \text{Cov}_{ij}(\ell) \) is the covariance of the spectrum \( \mathcal{C}_{\beta\gamma,i}(\ell) \) and is given by

\[ \text{Cov}_{ij}(\ell) = \frac{1}{2\ell + 1} \left[ \mathcal{C}_{\beta\gamma,i}(\ell) + \mathcal{C}_{\beta\gamma,i}(\ell) \right], \]

(35)

and where quantities with the bar represent the estimate of the signal, including intrinsic CMB fluctuations, instrumental noise and the beam of the CMB experiment as noise sources:

\[ \mathcal{C}_{\beta\gamma,i}(\ell) = C_{\beta\gamma,i}(\ell) \]

(36)

\[ \mathcal{C}_{\beta\gamma,i}(\ell) = C_{\beta\gamma,i}(\ell) + C_{\text{CMB}}(\ell) + w^{-1} B^{-2}(\ell) \]

(37)

4.1 iSW Fisher matrix

The Fisher matrix of the iSW effect follows directly from the variance of the spectrum estimates,

\[ F_{a,b}^{\text{iSW}} = \sum_{\ell} \frac{\partial \mathcal{C}_{\beta\gamma,i}(\ell)}{\partial \theta_a} \text{Cov}_{ij}(\ell) \frac{\partial \mathcal{C}_{\beta\gamma,i}(\ell)}{\partial \theta_b}. \]

(34)

4.2 Weak lensing Fisher matrix

The Fisher matrix for weak lensing is given by

\[ F_{a,b}^{\text{WL}} = f_{\text{sky}} \sum_{\ell} \frac{\partial \mathcal{C}_{\gamma\gamma,i}(\ell)}{\partial \theta_a} \text{Cov}_{ij}(\ell) \frac{\partial \mathcal{C}_{\gamma\gamma,i}(\ell)}{\partial \theta_b} \mathcal{C}_{\gamma\gamma,i}(\ell), \]

(39)

where the sum runs from \( \ell = 5 \) to \( 5000, \) (as from the official \textit{Euclid} prescriptions, see Laureijs et al. 2011), and where the sum over repeated indices is implied. We added a Poissonian shape noise term to the weak lensing spectra,

\[ \mathcal{C}_{\gamma\gamma,i}(\ell) = C_{\gamma\gamma,i}(\ell) + \delta_{ij} \frac{\gamma_{\text{int}}^{1/2}(\ell)}{n_i}, \]

(40)

\( \gamma_{\text{int}} \) is the rms intrinsic shear (here, we assume \( \gamma_{\text{int}}^{1/2} = 0.22 \)) and \( n_i \) is the number of galaxies per steradian belonging to the \( i \)th bin. We assume a Gaussian shape of the covariance while noting that non-Gaussian contribution can have a strong influence on the derived forecasts (Takada & Jain 2009).

4.3 Spectroscopic galaxy distribution Fisher matrix

The galaxy power spectrum Fisher matrix is given by (Seo & Eisenstein 2003)

\[ F_{a,b}^{\text{GC}} = \int_{k_{\text{min}}}^{k_{\text{max}}} \frac{k^2 dk}{4\pi^2} \frac{\partial \ln P_{\gamma\gamma}^{\text{spec}}(z; k, \mu)}{\partial \theta_a} \frac{\partial \ln P_{\gamma\gamma}^{\text{spec}}(z; k, \mu)}{\partial \theta_b} \times V_{\text{eff}}, \]

(41)

where GC stands for galaxy clustering, the observed galaxy power spectrum \( P_{\gamma\gamma}^{\text{spec}} \) is given by equation (33), the derivatives are evaluated at the parameter values of the fiducial model. \( k_{\text{min}} = 0.001 \) and \( k_{\text{max}} \) is such that the rms amplitude of the fluctuations at the corresponding scale \( R_{\text{max}} = 2\pi/k_{\text{max}} \) is \( \sigma(R_{\text{max}}) = 0.25, \) with an additional cut at \( k_{\text{max}} = 0.2 \) h/Mpc, in order to remain in the linear regime. \( V_{\text{eff}} \) is the effective volume of the survey, given by

\[ V_{\text{eff}} \simeq \left( \frac{n P_{\gamma\gamma}^{\text{spec}}(k; \mu)}{h P_{\gamma\gamma}^{\text{spec}}(k; \mu) + 1} \right)^2 V_{\text{survey}}. \]

(42)

The latter equation holds for an average comoving number density \( \bar{n} \). The number densities and further fiducial \textit{Euclid} specifications can be found in Laureijs et al. (2011) and Majorotto et al. (2012).

4.4 Forecasts

We computed forecasts on the measurement of \( c_\beta^2 \) and \( c_\gamma^2 \) for a wide range of fiducial values, in order to capture the parameter determining capability of experiments for a previously unknown set of
parameters. The probes are assumed to be uncorrelated as discussed above, hence their Fisher matrices add,

\[ F_{\alpha \beta} = F_{\alpha \beta}^{GC} + F_{\alpha \beta}^{WL} + F_{\alpha \beta}^{iSW}, \]

and we derive confidence contours on \( c_s^2 \) and \( c_v^2 \) and individual errors from this combined Fisher matrix, marginalizing over all other five parameters considered in this analysis. \( F_{\alpha \beta}^{GC} \) has been further marginalized over \( P_{\text{shot}} \), while the galaxy bias has been kept fixed to the fiducial values of Orsi et al. (2010) (see Section 3.3).

Our forecasts on the following fiducial models: \{\( c_s^2, c_v^2 \)\} = \{1, 0\}, \{10^{-3}, 10^{-4}\}, \{10^{-5}, 10^{-6}\} and \{10^{-6}, 10^{-6}\} are shown in Figs 3, 4, 5 and 6, respectively. The first model corresponds to the case of simple scalar field dark energy, while the following two pairs of fiducial models were chosen such that \( c_v^2 = 10^3 c_s^2 \). This is because, as mentioned previously in Section 2.2, the relevant quantity is the effective sound speed, and the effect of \( c_v^2 \) in it is \( \sim 10 \) times stronger than that of \( c_s^2 \) when \( w = -0.8 \) because of the factor multiplying \( c_s^2 \) in equation (14). The last model, also having small \( c_s^2 \) and \( c_v^2 \), does not verify the latter relation, and has been chosen in order to be compared to previous work (Sapone et al. 2013) and to the similar case \{\( c_s^2, c_v^2 \)\} = \{10^{-5}, 10^{-6}\}.

In all plots, iSW constraints are shown in blue, weak lensing ones in dark blue, GC ones in green, combined iSW-GC ones in yellow, combined iSW-weak lensing ones in orange, and combined iSW-GC-weak lensing ones in red.

From Figs 3–6 it is clear that the results depend very much on the chosen fiducial model. A common feature is that iSW on its own does not provide very strong constraints. In particular (see also Table 1), \( c_v^2 \) is quite badly constrained, with relative errors ranging between \( 6.1 \times 10^4 \) and \( 1.1 \times 10^5 \), while relative errors on \( c_s^2 \) are much smaller: between \( 1.1 \times 10^{-1} \) and \( 1.4 \times 10 \). This was to be expected since the iSW effect alone has a rather small signal strength of about \( 5\sigma \) for cross-correlating the CMB with the Euclid galaxy sample (Douspis et al. 2008).
Weak lensing constraints\textsuperscript{4} are much stronger than iSW ones in the case of a fiducial scalar field dark energy, but become progressively comparable to them when the fiducial $c_s^2$ and $c_r^2$ become smaller, with the exception of the case $c_s^2 = c_r^2 = 10^{-6}$.

Even though both iSW and weak lensing do not give very strong constraints on sound speed and viscosity (see also Table 1), it is very interesting to notice that the two data sets complement each other very well. This is especially true for the case $c_s^2 = 10^{-5}$ and $c_r^2 = 10^{-6}$, as represented in Fig. 5, where the blue ellipses, which indicate errors from iSW, have a very different degeneration direction with respect to the dark blue contours, corresponding to errors from weak lensing, but are comparable to them in size. Therefore the resulting combined errors are much smaller than those from a single data set. In particular, the iSW effect gives better constraints on the sound speed and weak lensing on the viscosity parameter. Also in the cases of Figs 3, 4 and 6 there is a visible but small improvement when adding iSW to the other data sets. The reason for this is that iSW’s sensitivity to $c_r^2$ and $c_s^2$ is comparable to that of GC and WL, but its sensitivity to all other parameters is much less than that of the other observables. This can be seen by comparing the Fisher matrix elements, and it is demonstrated in the example of Fig. 7. Here we show how the error on $c_r^2$ from GC, for the fiducial model with $c_r^2 = 10^{-6}$ and $c_s^2 = 10^{-6}$, is modified when adding to the GC Fisher matrix progressively larger iSW sub-matrices. The first bar on the left corresponds to $\sigma_{c_s^2}^2$ from GC only, while the last one on the right corresponds to $\sigma_{c_r^2}^2$ from joint GC-iSW data. The second bar from the left corresponds to adding to GC only the $c_r^2 - c_s^2$ element of the iSW Fisher, the third to adding the $c_r^2 - \Omega_m h^2$ sub-matrix, and so forth.

\textsuperscript{4}With respect to Sapone et al. (2013) we have improved the estimation of $P_{NL}$ by using the full \texttt{CAMB} output instead of an analytical approximation to it.

Figure 7. Errors on $c_r^2$, for the fiducial model with $c_r^2 = 10^{-6}$ and $c_s^2 = 10^{-6}$, from the GC Fisher matrix joined with progressively larger iSW sub-matrices. The first bar on the left corresponds to $\sigma_{c_s^2}^2$ from GC only, while the last one on the right corresponds to $\sigma_{c_r^2}^2$ from joint GC-iSW data. The second bar from the left corresponds to adding to GC only the $c_r^2 - c_s^2$ element of the iSW Fisher, the third to adding the $c_r^2 - \Omega_m h^2$ sub-matrix, and so forth (see legend). As can be seen, if iSW were depending only on $c_s^2$ and were independent of all other parameters, combining iSW with GC would improve the error on $c_r^2$ by a factor of about 3 (see second bar from the left). Since iSW also depends on other parameters, which are correlated to $c_r^2$, the improvement becomes weaker and weaker the more parameters are added. Hence, when joining the data sets there is a definite improvement in measurements of these two parameters, due to the improved sensitivity to $c_s^2$ and $c_r^2$ and to the better sensitivity to all other parameters due to GC and WL. Instead, the errors on $c_r^2$ and $c_s^2$ from iSW alone become quite large after marginalizing over all other parameters, which are correlated to the sound speed and the viscosity parameter.

Also errors on $c_s^2$ and $c_r^2$ from GC are orthogonal to those from the iSW, but only in the case where $c_s^2 = 10^{-5}$ and $c_r^2 = 10^{-6}$ (see Fig. 5) this helps reducing the errors, because the former data set performs better in constraining $c_s^2$ and the second $c_r^2$. In the case of fiducial $c_s^2 = 10^{-5}$ and $c_r^2 = 10^{-3}$ it is GC which gives best errors on both parameters, while for the $c_s^2 = 1$ and $c_r^2 = 0$ fiducial model, it is weak lensing.

Another interesting question is whether iSW adds important information to that provided by the other two data sets, which had already been analysed in Sapone et al. (2013). Table 1 answers this
question. It turns out that the information from iSW helps significantly in constraining $c^2_s$ and $c^2_v$ if the true model has $c^2_s = 10^{-5}$ and $c^2_v = 10^{-6}$, see Fig. 5. In this case, the iSW alone gives a strong constraint, which has moreover a different degeneration direction with respect to the error from galaxy clustering and weak lensing. For the other three fiducial models the gain when adding iSW is not very strong, as in both cases the combination of weak lensing and galaxy clustering gives a much tighter constraint in both the sound speed and the viscosity than the iSW alone.

It is also interesting to see at how many $\sigma$ one can detect a deviation from $c^2_s = 0$ or from $c^2_v = 0$. This is shown in Table 2. In particular, already with iSW alone one can have a $7.3\sigma$ detection of $c^2_s \neq 0$ for the fiducial $[c^2_s, c^2_v] = [10^{-5}, 10^{-6}]$. To have a detection of more than $3\sigma$ of $c^2_s \neq 0$ one has to combine iSW with WL or GC, where the latter observable gives a stronger detection. When combining all observables, one obtains the best detection for the fiducial $[c^2_s, c^2_v] = [10^{-5}, 10^{-6}]$: for the case $[c^2_s, c^2_v] = [10^{-5}, 10^{-6}]$, one has a strong detection of $c^2_s \neq 0$ but only a $1.7\sigma$ detection of $c^2_v \neq 0$. In the fiducial $[c^2_s, c^2_v] = [10^{-5}, 10^{-6}]$ instead one detects better (at $2.9\sigma$) $c^2_v \neq 0$ while $c^2_s \neq 0$ keeps being undetected.

It is interesting to notice (see Fig. 6) that when both the sound speed and the viscosity are small, but the relation between $c^2_s$ and $c^2_v$ differs from $c^2_v \sim 10^{-2}$, the iSW effect error ellipse becomes much larger and as a result the sound speed parameter is less strongly constrained. Thus, we conclude that (i) very interesting results can be obtained through a combination from different cosmological probes and that (ii) the iSW effect is able to tighten constraints significantly for cases where there is a large difference between $c^2_s$ and $c^2_v$.

### 5 SUMMARY

In this paper, we have investigated how well the viscosity and sound speed of dark energy can be measured with the iSW cross-correlation spectrum, when using Planck and Euclid observations, and how joining iSW measurements to galaxy clustering and weak lensing ones improves constraints.

We found that the speed of sound is quite well constrained, with relative errors as small as 0.14 for small fiducial $c^2_s$ and $c^2_v$, while relative errors on the viscosity parameter are very large. The anisotropic stress is not well constrained by the iSW, due to the marginalization over the other parameters, which degrades the sensitivity to the sound speed and viscosity. In practice, iSW is only sensitive to $c^2_s$ and $c^2_v$ and not to the other parameters as it is for GC and WL. However, the error ellipses are interestingly orthogonal to those from weak lensing and the latter has much better sensitivity to all the other parameters. Hence, the combination of these two data sets constrains tightly the parameter space, giving relative errors on $c^2_s$ and $c^2_v$ as small as $2.6 \times 10^{-2}$ and $7.5 \times 10^{-2}$, respectively. This is an improvement of a factor $\sim 1500$ in the measurement of the sound speed and $\sim 50$ in the measurement of the viscosity parameter, with respect to the weak lensing only constraint. The improvement obtained when combining iSW with galaxy clustering is smaller: a factor of $\sim 1.5$ in $c^2_s$ and $\sim 150$ in $c^2_v$. Finally, the addition of iSW to weak lensing and galaxy clustering constraints is important if the fiducial sound speed and viscosity parameter are very small, while it is not very relevant for higher fiducial values of $c^2_v$.

It is also important to remind that in order to make the effect of dark energy perturbations stronger we have always used a value of the equation of state parameter $w = -0.8$. For values close to $w = -1$ the effects on the observables due to the dark energy perturbations are reduced, as all the phenomenological functions used (such as $Q(k, a)$) have a term $\propto (1 + w)$. If we use a value of $w = -0.9$ we expect our final errors on the parameters to increase. But by how much? All the observables used in this paper depend most strongly on $Q^2$ (see equation 15) which is intrinsically included into the matter power spectrum; for a sound speed equal to zero $Q = 1 = (1 + w)/(1 - 3w)a^{-3w}$ so the relative increase of the errors on the sound speed will be given by $1/[(Q(w = -0.9) - 1)/(Q(w = -0.8) - 1)]^2$ which is of about a factor 4 larger, in agreement also with the results found in Sapone, Kunz & Amendola (2010).

A detection of sound speed and viscosity different from the values associated with a classical scalar field, i.e. $c^2_s = 1$ and $c^2_v = 0$, will point to a new understanding of the accelerated phase of the Universe. This is because the non-ideal fluid considered in this paper can be thought of as an effective dark energy fluid parametrizing a modified gravity model, see Kunz & Sapone (2007). In practice, the detection of a zero sound speed does not automatically mean that we are dealing with an actual dark energy fluid, even though one would nevertheless experience effects which could be attributed to fluctuations of a fluid.

In this paper, we found that joining data from Euclid and Planck we are able to constrain simultaneously the sound speed and the viscosity parameters, provided that the two are sufficiently small. This is mostly due to the different sensitivity of the three observables, i.e. GC, WL and iSW to the two parameters. In most cases, the iSW has a different degeneracy with respect to WL and GC, and this helps reducing the errors on $c^2_s$ and $c^2_v$ by a factor of $\sim 100$ (as pointed out before). Our results are in agreement with what found by Mota et al. (2007), Calabrese et al. (2011), Chang & Xu (2014), who show that for values of $c^2_v$ approaching $c_s^2 = 1$ the detection of a positive viscosity is very difficult, even when, as in the case of Calabrese et al. (2011), an early dark energy helps its detection by increasing its effect on smaller scales.

To conclude, are Euclid and Planck able to measure the sound speed and the viscosity parameters of a dark energy component?
If the values of $c_s^2$ and $c_v^2$ are small enough, the answer is yes; consequently, we will be able to constrain well the effective dark energy model. On the contrary, if sound speed and viscosity will escape detection, at least one of the two parameters will likely have large values. We would assume that other cosmological probes would not directly provide constraints on dark energy properties, but would nevertheless be able to provide constraining power by fixing other parts of the cosmological model, such as the dark matter density or the dark energy equation of state, which was not subject to variation in our investigation.

ACKNOWLEDGEMENTS

We acknowledge Martin Kunz and Luca Amendola for inspiring discussions. EM was supported by the Spanish MINECO's 'Centro de Excelencia Severo Ochoa'—programme under grant no. SEV-2012-0249 and by the Spanish MICINN's Juan de la Cierva programme. EM was supported by the Spanish MINECO’s ‘Centro de Excelencia Severo Ochoa’—programme under grant no. SEV-2012-0249 and by the Spanish MICINNs Juan de la Cierva programme (JCI-2010-08112), by CICYT through the project FPA-2012-31880, by the Madrid Regional Government (CAM) through the project HEPHACOS S2009/ESP-1473 under grant P-ESP-00346 and by the European Union FP7 ITN INVISIBLES (Marie Curie Actions, PITN-GA-2011-289442). DS acknowledges financial support from the Fondecyt project number 11104096 and from the 'Anillo' project ACT1122 founded by the 'Programa de Investigación asociativa'.

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