Measuring Interstellar Delays of PSR J0613−0200 over 7 years, using the Large European Array for Pulsars

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ABSTRACT

Using data from the Large European Array for Pulsars (LEAP), and the Effelsberg telescope, we study the scintillation parameters of the millisecond pulsar PSR J0613−0200 over a 7 year timespan. The “secondary spectrum” – the 2D power spectrum of scintillation – presents the scattered power as a function of time delay, and contains the relative velocities of the pulsar, observer, and scattering material. We detect a persistent parabolic scintillation arc, suggesting scattering is dominated by a thin, anisotropic region. The scattering is poorly described by a simple exponential tail, with excess power at high delays; we measure significant, detectable scattered power at times out to $\sim 5\, \mu s$, and measure the bulk scattering delay to be between 50 to 200 ns with particularly strong scattering throughout 2013. These delays are too small to detect a change of the pulse profile shape, yet they would change the times-of-arrival as measured through pulsar timing. The arc curvature varies annually, and is well fit by a one-dimensional scattering screen $\sim 40\%$ of the way towards the pulsar, with a changing orientation during the increased scattering in 2013. Effects of uncorrected scattering will introduce time delays correlated over time in individual pulsars, and may need to be considered in gravitational wave analyses. Pulsar timing programs would benefit from simultaneously recording in a way that scintillation can be resolved, in order to monitor the variable time delays caused by multipath propagation.

Key words: pulsars: general – pulsars:individual ( PSR J0613−0200 )

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1 INTRODUCTION
Radio emission from pulsars experiences several propagation effects from the ionised interstellar medium (ISM), as the index of refraction varies with electron density and frequency. The signal acquires a group delay $\Delta t$, known as dispersion, scaling as $\Delta t \propto DM v^{-2}$, where $DM$ is the integrated column density of free electrons, and $v$ is the observing frequency. Spatial variations in the electron density result in multi-path propagation, with deflected paths acquiring a geometric time delay from the path-length difference compared to the direct line-of-sight. When these delays are large (compared to the pulse duration), it is observed as scintillation, the one-sided broadening of pulses often resembling an exponential tail. When these delays are small, we observe it as scintillation, the constructive and destructive interference of different deflected im-

tages of the pulsar (stationary phase points, regions where light can be deflected to the observer). Each image has a geometric time delay $\tau_i$ and a fringe rate (or Doppler rate) $f_{D,i}$, with a magnification $\mu_i$ and intrinsic phase $\phi_i$. In this approximation, the contribution of all of the images is

$$g_E(\tau, f_D) = \sum_i \sqrt{\mu_i} e^{-i\phi_i} \delta(f_D - f_{D,i}) \delta(\tau - \tau_i).$$  (1)

One of the central goals of pulsar timing is to directly detect gravitational waves, in a so-called pulsar timing array (Hobbs, Desvignes et al. 2016, Verbiest et al. 2016, Arzoumanian et al. 2018). The most stable pulsars are observed on weekly to monthly cadence over many years, and $\sim$nHz gravitational waves could be observed in timing residuals correlated in time and position on the sky (Hellings & Downs 1983). This effect is expected to be tiny, with a fractional change of the arrival time compared to the gravitational wavelength of order $10^{-15}$, so it requires careful understanding of all other effects which would change the arrival times of pulses. While PTA pulsars are selected for their stability, they all experience variable dispersion and scattering to some degree due to the relative motion of the pulsar and observer with respect to the ISM. Variable dispersion measures have been measured and corrected using multifrequency data (eg. Keith et al. 2013), while changes in scattering time are often estimated using the statistical relation between the scintillation bandwidth (the frequency width of scintillation) and scattering time (eg. Levin et al. 2016; Shapiro-Albert et al. 2019, see Verbiest & Shaifullah (2018) for a review of how these effects limit precision pulsar timing). Dispersion and scattering both scale strongly with frequency, and are often covariant. One can look for variable delays following a $v^{-2}$ scaling (Lam et al. 2019); the technique of wide-band template matching has recently been developed as a way to jointly fit for these effects (Liu et al. 2014; Pennucci et al. 2014; Pennucci 2019; Alam et al. 2020).

In this paper, we begin to apply the methods of Hemberger & Stinebring (2008) to PTA pulsars, in which scintillation arcs are used to estimate time delays from multi-path propagation. We analyse PSR J0613–0200 over 7 years, in roughly monthly cadence, using data from the Large European Array for Pulsars (LEAP) (Stappers & Kramer 2011; Bassa et al. 2016), and a 3-month bi-weekly observing campaign using the 100–m Effelsberg radio telescope. This pulsar is of particular interest; it shows the strongest evidence of a 15 nHz strain, but since the signal appears most strongly in this pulsar, it is believed to arise from an unmodelled non-GW signal (Aggarwal et al. 2019). In Section 2, we give an overview of some necessary background of scintillation, and summarise the methods of Hemberger & Stinebring (2008). In Section 3, we describe our observations with the LEAP telescope, our short-term observing campaign with the Effelsberg telescope. In Section 4 we outline our methods, in Section 5 we present our results, and we discuss the ramifications and future prospects in Section 6.

2 BACKGROUND ON THEORY OF SCINTILLATION

2.1 Thin screen theory and stationary phase approximation
The theory of scattering in thin screens is outlined in detail in Walker et al. (2004) and Cordes et al. (2006), and we summarise some of the pertinent relations here.

The “stationary phase approximation” assumes that the observed signal can be described as a coherent summation over all images of the pulsar (stationary phase points, regions where light can be deflected to the observer). Each image has a geometric time delay $\tau_i$ and a fringe rate (or Doppler rate) $f_{D,i}$, with a magnification $\mu_i$ and intrinsic phase $\phi_i$. In this approximation, the contribution of all of the images is

$$|I(\tau, f_D)|^2 \approx \sum_{ij} \delta(f_D - f_{D,ij}) \delta(\tau - \tau_i) \delta(\tau + \tau_j),$$  (2)

where $f_{D,ij}$ and $\tau_{ij}$ are the differences between two interfering images,

$$f_{D,ij} = \frac{(\theta_i - \theta_j) \cdot v_{\text{eff}}}{\lambda},$$  (3)

$$\tau_{ij} = \frac{d_{\text{eff}}(\theta_i^2 - \theta_j^2)}{2c}.$$  (4)

We note that, since the dynamic spectrum is a real function, the secondary spectrum is point-symmetric.

The effective distance $d_{\text{eff}}$ and effective velocity $v_{\text{eff}}$ depend on the fractional distance of the screen from the pulsar $s$ as

$$d_{\text{eff}} = (1/s - 1) d_{\text{par}},$$  (5)

$$v_{\text{eff}} = (1/s - 1) v_{\text{par}} + \frac{v_\parallel - v_{\text{act}}}{s},$$  (6)

$$s = 1 - d_{\text{act}}/d_{\text{par}}.$$  (7)

Where, $d_{\text{par}}$ and $v_{\text{par}}$ are the pulsar’s distance and velocity, $v_\parallel$ and $v_{\text{act}}$ are the velocities of the Earth and scattering screen, respectively (and where we are only considering the 2D velocity on the plane of the sky).

Considering one image as the direct line-of-sight, then $\theta_i = \theta_j = 0$, and $\tau$ and $f_D$ are related through their common dependence on $\theta$:

$$\tau = \eta f_D^2, \quad \text{with} \quad \eta = \frac{d_{\text{eff}} \lambda^2}{2c v_{\text{eff}}^2},$$  (8)

where $\lambda$ is the observing wavelength, and where $v_{\text{eff}}$ is the effective velocity along the position vector $\theta$ to the image. For a 1D distribution of images, we denote the angle of the screen’s axis with $\psi$. Many images along a line interfering with the direct line of sight then results in a parabolic distribution of power in the secondary

\footnote{1 equivalently, one must Fourier transform over a long enough timespan of $E(t)$ fully encompassing $g(t)$ by a factor of 2, due to the Nyquist Theorem – the longest timescales correspond to the finest frequencies}
spectrum, while the commonly seen “inverted arclets” (eg. Stinebring et al. 2001) arise from the interference between subimages. The curvature $\eta$ depends on the distance to the screen, the effective velocity, and the angle between the velocity and the screen. Structures in the secondary spectrum move along the main parabola from left to right (negative to positive $f_D$) due to the effective velocity as

$$\frac{df_D}{df} = \frac{1}{2\eta^V} \left(1 - \nu f_D \frac{d\eta}{df}\right). \quad (9)$$

The motion of points in the secondary spectrum is uniquely defined by the curvature of the parabolic arc and its time-derivative – in other words, clumps of power in the secondary spectrum must move, and the resulting bulk scattering time is necessarily variable. Variable motion from the Earth’s or the pulsar’s orbit will contribute to $\nu f_D \frac{d\eta}{df}$.

### 2.2 The interstellar response

In this section we summarise and expand upon the method of Hemberger & Stinebring (2008), to use the secondary spectrum to estimate the total time delays from multipath propagation.

The electric field that we observe is the intrinsic signal of the pulsar convolved with the impulse response function of the ISM,

$$E(t) = (E_{\text{int}} * g_E)(t), \quad (10)$$

where $E_{\text{int}}$ is the intrinsic signal of the pulsar, and $g_E(t)$ is the interstellar impulse response function of the field.

We measure the time-averaged intensity, not the direct electric field. The quantity of interest is then the time shift of the intensity $\langle \tau \rangle_{I(t)}$, where

$$I(t) = \langle |E(t)|^2 \rangle = \langle (|E_{\text{int}} * g_E)(t)|^2 \rangle, \quad (11)$$

where $\langle \rangle$ denotes the average over many pulses. First we must find a suitable way to describe the effect of response of the field $g_E(t)$ on the intensity. Under the assumption that the intrinsic field is temporally incoherent, then

$$\langle E_{\text{int}}(t_1) E_{\text{int}}(t_2) \rangle = \delta(t_1 - t_2), \quad (12)$$

and it can be shown that the observed intensity can be written as

$$I(t) = (I_{\text{int}} * g)(t) \quad (13)$$

where $g(t) = |g_E(t)|^2$ can be thought of as the intensity response function. Equation 12 is written unrigorously for infinite bandwidth – in the real case of a finite bandwidth, the delta function would be replaced by a sinc function with width $\sim 1/BW$.

With the above assumptions, we now have the intensity written in a form resembling Hemberger & Stinebring (2008), and can follow their steps. The goal of this method is to estimate the time shift from the intensity response function, $\langle \tau \rangle_{I(t)} \equiv \tau_i$. Two properties of convolutions are important for this method, that the centroid and the variance of two convolved functions are additive

$$\langle \tau \rangle_{f_g} = \langle \tau \rangle_{f} + \langle \tau \rangle_{g}, \quad \text{and} \quad \sigma_{f_g}^2 = \sigma_f^2 + \sigma_g^2 \quad (14)$$

For simplicity, we define the center of the pulse to be at $t = 0$, and define the pulse width to be $w$. When $w \gg \tau_i$, as is the case for this paper, as the $\sim \mu$s delays are much smaller than the $\sim$ms pulse width), then using the two convolution properties above, we have

$$\langle \tau \rangle_{I} = \langle \tau \rangle_{I_{\text{int}}} + \langle \tau \rangle_{g_E} = \tau_i, \quad (15)$$

and

$$\sigma_i^2 = \sigma_{I_{\text{int}}}^2 + \sigma_{g_E}^2 \approx w^2 + \tau_i^2 \approx w^2, \quad (16)$$

since $\langle \tau \rangle_{I_{\text{int}}} = 0$ by definition. This means that the shape of the pulse is effectively unchanged, yet it still has a bulk time delay from the response.

### 2.3 Estimating time delays from the Secondary Spectrum

Now we address how to estimate the time delays of the intensity response function in practice. As we are only concerned with measuring time delays, in this section we drop the time dependence for simplicity, focusing on the imprint of the impulse response function on the spectrum.

The Fourier transform of the intensity spectrum is

$$I(\nu) = \tilde{I}_{\text{int}}(\nu) \tilde{g}(\nu), \quad (17)$$

where the convolution between the pulsar’s signal and impulse response becomes a direct multiplication. The intrinsic profile $\tilde{I}_{\text{int}}$ is assumed to be stable, and only slowly varying across frequency after averaging over many pulse rotations, so we treat it as a constant. The secondary spectrum is obtained by Fourier transforming and squaring the spectrum $I(\nu)$, resulting in

$$|I(\nu)|^2 = \tilde{I}_{\text{int}}(\nu) |g(\nu - \nu)| \tilde{g}'(\nu). \quad (18)$$

This is the autocorrelation of the intensity impulse response function, and we see that this form cannot necessarily recover the total time delay as it only measures differences in $\tau$, not an absolute time.

To simplify, we return to the stationary phase approximation, as discussed in Section 2.1. The intensity response function is the square modulus of the field as given in equation 1, where if we assume that the images lose coherence when integrating over the full observation we have

$$g_i(\tau) = \sum \mu_i \delta(\tau_i). \quad (19)$$

We wish to estimate this from the secondary spectrum. To examine a limiting case, let us assume most of the power is near the undeflected line of sight (defined as $j = 0$), then $\tau_0 = 0$, $\mu_0 \approx 1$, and $\mu_0 \gg \mu_i$. Then, taking only positive $\tau_i$, and averaging over $f_D$, equation (2) becomes

$$|I(\tau)|^2 \approx \sum_i \mu_i \delta(\tau_i) \quad (20)$$

In this limit, there will be a visibly strong parabolic arc without inverted arclets. The total time delay would then be determined from the expectation value in $\tau$, where the contribution of the bright central image divides out

$$\langle \tau \rangle_{I} = \sum_i \mu_i \tau_i \quad (21)$$

$$= \sum_i \mu_i \tau_i \quad (22)$$

$$= \langle \tau \rangle_{g} \quad (23)$$

The contribution of the phases can be neglected if every pixel in the secondary spectrum contains only one pair of interfering images – while not necessarily the case, this is aided by the time axis of the dynamic spectrum and many channels, which separates the power in the secondary spectrum in $f_D$ as well as $\tau$.

We see that in the limit of a strong central image, we can recover the total time delays from the secondary spectrum. More
generically, how well the time delays can be computed from the secondary spectrum is dependent on the unknown distribution of images (or the functional form of $g_i(\tau)$). In the case of strong scattering, there is no reason to expect a single undeflected line of sight image, but rather there may be many bright, scattered images at small angular separation. In this case, the time delay computed from the above formula will be overestimated, due to the cross-terms of bright central images interfering. This will bias the result high by a factor of $\sim 2m/(m+1)$, where $m$ is the number of bright images, leading to a difference as large as a factor of 2. In the case of a discretised secondary spectrum, this will only begin to matter if the image separations are larger than one pixel in $\tau$, otherwise it will approximate the case of one bright central image. Additionally, we are still limited by the fact that the secondary spectrum measures differences in time delays, rather than absolute time delays; if there is a time-shift applied to all images, it would not be captured by our estimate.

With the above caveats mentioned, we use equations 20 and 23 as our basis to measure time delays throughout the paper. These include the assumption of a strong central image, which we believe gives a reasonable estimate for our purposes. We describe how to compute time delays in practice from our data in Section 4.3, after detailing our data reduction and secondary spectra creation.

3 OBSERVATIONS

3.1 LEAP

The Large European Array for Pulsars (LEAP) is a phased array of five large radio telescopes in Europe; the Effelsberg telescope, the Lovell telescope at Jodrell Bank Observatory, the Westerbork Synthesis Radio Telescope, the Nançay Radio Telescope and the Sardinia Radio Telescope (Stappers & Kramer 2011). The coherent addition of radio signals from all these telescopes results in an effective 195 m diameter dish. The overview of LEAP is given in Bassa et al. (2016). Observations have been made monthly since 2012, with whichever subset of these telescopes was available.

The voltage data from each site are shipped or transferred to Jodrell Bank Observatory to be correlated and coherently added on a designated CPU cluster, using a specifically designed software correlator (details in Smits et al. 2017). Correlation involves a polarisation calibration based on an observation of PSR J1022+1001 or PSR B1933+16 from the same epoch, correlation on a calibrator to find an initial phasing solution, then self-calibration on the pulsar to determine the time delays and fringe drift rates for each telescope throughout the observation, using Effelsberg as a time and position reference. The coherently added voltages are stored on tape, allowing us to re-reduce the data with arbitrary time or frequency resolution. The high sensitivity, and the flexibility offered by storing the baseband data has enabled LEAP to do single pulse studies of MSPs (Liu et al. 2016; McKee et al. 2019); for these same reasons, it is an ideal telescope for the scintillation work presented in this paper. Typical observing lengths are 30 – 60 minutes, with bandwidths of 80 – 128 MHz (comprised of 16 MHz subbands), depending on the subset of telescopes used for a given observation. As we will show in Section 5.2, the angular extent of the scattering screen is unresolved by LEAP, so we can safely treat it as a single-dish instrument for our purposes.

3.2 Effelsberg 100-m Telescope

From March to June 2020, we had a roughly bi-weekly monitoring campaign using the Effelsberg telescope. Baseband data were recorded as 8-bit “dada” files using the PSRIX backend (described in Lazarus et al. 2016), using the central feed of the 7-beam receiver (“P217mm”). The data were recorded in 25 MHz subbands, with a usable bandwidth of 1250-1450 MHz, and typical observation lengths of 90 minutes. While Effelsberg alone is less sensitive than LEAP, this is compensated through the larger exposure times and bandwidth.

4 METHODS

4.1 Creating dynamic and secondary spectra

We created folded archives from the baseband data using dapar (van Straten & Bailes 2011), coherently de-dispersing and folding with 10 s bins, 128 phase bins, and sufficient channels to fully resolve scintillation - 62.5 kHz, and 50.0 kHz channels for LEAP and Effelsberg respectively. The subbands were combined in frequency using the parcfile tool paradd (Hotan et al. 2004) to form one combined archive per observation. The following processing steps for data from either telescope are identical unless expressly stated otherwise.

After summing polarisations, each folded archive contains a data cube $F(t, \nu, \phi)$. We use a fixed off-pulse region relative to the pulse, a contiguous 50% section with no apparent pulsed emission (in Figure 2, phase 0.5–1.0) We divide by the time average of the off-pulse region across the full observation to approximately remove the bandpass, and in each time and frequency bin, we subtract the mean of the off-pulse region to remove variable background flux. Sub-integrations with an off-pulse standard deviation $> 5\times$ the mean rms value were masked, as were any time bins or frequency channels with $> 30\%$ of flagged sub-integrations.

The LEAP dada files are saved separately in each sub-band, in individual 10 s files; a small number of these files were missing, and were filled with zeros, and included in our mask. To reduce artefacts caused by Fourier transforming over a window function, masked pixels were iteratively in-painted using the mean of the nearest pixels. While more sophisticated methods of inpainting exist, this is sufficient for our analysis, as typically no more than 5% of data are flagged.

A time and frequency averaged profile was created, and zeroed everywhere the S/N was below 5 $\sigma$. This profile was used to weight each phase bin, before summing over pulse phase to create the dynamic spectrum $I(t, \nu)$. Over a narrow band, it is sufficient to simply use a 2D FFT, which we used for this analysis (over a wider band, the $\nu^{-2}$ scaling of $\eta$ causes arcs to smear in the secondary spectrum, summarised in Gwinn & Sosenko 2019). Before taking a FFT, we padded the edges by a factor of two with the mean value of the central bins. This profile was used to form the baseband for the scintillation work presented in this paper. Typical observing lengths are 30 – 60 minutes, with bandwidths of 80 – 128 MHz (comprised of 16 MHz subbands), depending on the subset of telescopes used for a given observation.

4.2 Measuring arc curvatures

The main power in the secondary spectrum of PSR J0613–0200 follows a parabolic arc, suggesting scattering dominated by a

http://psrdada.sourceforge.net/
Interstellar delays of J0613-0200 with LEAP

Figure 1. Top: Dynamic Spectra of 5 observations around the same time of year, to have comparable contributions from the Earth’s velocity. The colourbar extends from 2σ below the mean to 5σ above. Bottom: Corresponding secondary spectra, with a logarithmic colourbar extending three orders of magnitude. Clear arcs with noticeable localized clumps of power are seen, these correspond to prominent diagonal features in the above dynamic spectra. The observation from 2013 is anomalous, showing extremely fine stripes in the dynamic spectrum, corresponding to power at large time delays.

4.3 Integrating the secondary spectrum

For purposes of measuring time delays, the x-axis $f_D$ is not important, except to localise the scattered power in this parameter space. We isolate the power in a 1mHz region surrounding the main arc, as defined by our measured arc curvatures. We subtract the averaged background far from the main arc, assuming the noise is well described as a function of time-delay. We measure the total time delay through the expectation value of $\tau$, computed as

$$\langle \tau \rangle = \left|\frac{\int_{-T}^{T} |I(\tau)|^2 d\tau}{\int_{0}^{T} |I(\tau)|^2 d\tau}\right|,$$

(24)

where $T = 8\mu s$, defined by our choice of channelisation.

Artefacts in the dynamic spectrum, such as RFI, phasing imperfections, and the window function lead to correlated features in the secondary spectrum. As such, the noise properties are not always well behaved, and direct error propagation underestimates the error on $\langle \tau \rangle$. We estimate our errors directly from the cumulative function in equation (24); at high enough $T$ the integral plateaus, with residual variations caused by the effect of integrating noise in the secondary spectrum. We take the mean and standard deviation of equation (24) between $T = 4 - 8\mu s$ as our measurement and error of $\langle \tau \rangle$ respectively.
4.4 “Timing” a convolved template

To illustrate the effects of scattering on a profile, we can directly convolve our measure of the amplitude of $g_s(\tau)$ into a template profile and measure the time offset using the standard Fourier template-matching algorithm outlined in the appendix of Taylor (1992). We create an analytic template using the standard pauas tool (Hotan et al. 2004), fitting the profile with a series of von Mises functions, and interpolate the solution to have the equivalent 31.25 ns bins of our measured $I(\tau)^2$. We convolve the two, and measure the relative time delay between the convolved template against the original one. Figure 2 shows this convolution applied to one of our observations. The measured time delay in this way agrees perfectly with the method in the previous section; timing recovers the shift correctly, even when the effects are not visibly noticeable.

We note again that this is not precisely the intensity impulse response, but rather its autocorrelation, but it is close enough in amplitude to demonstrate that the convolved template is visually identical (with residuals at the 0.1% level after aligning the template), yet is measurably delayed. In addition, $I(\tau)/2$ is noticeably clumpy and poorly described by an exponential, even after being effectively smoothed by the autocorrelation.

4.5 Inferred time delay from the frequency ACF

A standard way to infer the time delays from scattering is to construct the auto-correlation functions $R(\Delta \nu) = (I(\nu)I(\nu))$. Fitting the width (specifically, the HWHM) of the ACF in frequency gives the scintillation bandwidth $\nu_{\text{scint}}$, which is inversely proportional to the bulk scattering delay as $\langle \tau \rangle = C/2\nu_{\text{scint}}$ (C is commonly assumed to be 1, and depends on the assumptions of the scattering distribution). This method is often used when arcs cannot be resolved nicely, as $\Delta \nu$ can typically be measured simply and robustly. However, if $g_s(\nu)$ is not smooth, then the ACF will be poorly described as a single Gaussian. Two such examples of an ACF, one well described, and one poorly described by a 1D Gaussian fit are shown in Figure 3, along with their derived $\nu_{\text{scint}}$ and $\langle \tau \rangle$. In the second case, a single Gaussian would preferentially fit the broad component, and result in an inferred time delay which is low, while power at large time delays results in the narrow peak smaller than 1 MHz.

Our error includes both the measurement error of the fit, in addition to the “finite scintle error”, which is a counting error of $\sqrt{N_{\text{scint}}}$, estimated in the same manner as Levin et al. (2016) as

$$\delta(\tau)/\langle \tau \rangle \approx (1 + \eta_1 T_{\text{obs}}/\nu_{\text{scint}})(1 + \eta_2 B/W/\nu_{\text{scint}})^{-1/2}.$$ (25)

The values of $\eta$ are the filling fraction of scintles, assumed here to be 0.2.
5 RESULTS

5.1 Evolution of scattering time from LEAP: month to year timescales

We present our measurements of the bulk time delay in the top panel of Figure 4. We find significant persistent scattering at the ~80 ns level, and a few cases of strong scattering variability on several month to year timescales. The timescales are set by the time it takes power to move through the secondary spectrum, as described in Section 2. The most striking feature is the strong excess scattering in 2013, where the bulk scattering is variable and extends above 200 ns. This is not captured very well by the ACF method, as the Gaussian fit latches onto the broad-scale scintillation rather than the narrow peak caused by the large time delays, as described in Section 4.5. As hinted at in Figure 3, this could potentially be remedied by using a multi-component model to the ACF, as in principle the ACF contains the equivalent information as the secondary spectrum, only differing by a Fourier transform.

In the bottom panel of Figure 4, we plot the DM values of PSR J0613−0200 from NANOGrav’s 12.5 year release (Alam et al. 2020a). The DM is steeply decreasing prior to 2013, and there is clear annual variation; this was studied in detail in Jones et al. (2017), with data spanning from 2006 until near the end of 2013. The authors fit the time variations of DM with a 1-year period sinusoid and a linear trend, to capture the contribution of the pulsar’s orbital motion. We find significant persistent scattering at the beginning of 2013. The time of closest approach for PSR J0613−0200 is in mid-June, and features can be seen to cross the line of Ω = 0, indicating a changing sign of vscr. To properly account for the arc curvature then depends only on the effective velocity parallel to the screen. By measuring the change in arc curvature over the year, one can measure the distance and orientation of the scattering screen.

We perform only a simple analysis here, currently ignoring the contribution from the pulsar’s orbital motion. We fit the observed curvature values beyond 2013 with a 1-dimensional screen, using measured values of the pulsar’s distance of 780 ± 80 pc and proper motion of μα = 1.822 ± 0.008 mas/yr, μδ = −10.355 ± 0.017 mas/yr from Desvignes et al. (2016). The three free parameters are the fractional screen distance s, the orientation of the screen Ψ, and vscr, the velocity of the scattering screen parallel to its axis of anisotropy. A 1D screen fits the data well, shown in the middle panel of Figure 4, while the best fit values are in Table 1. Using the screen distance, and the largest detectable time delays of τ ≈ 5 μs, the largest angular extent of the screen is θ ≈ 9°, smaller than the resolution of the longest baselines of LEAP.

During the increased scattering of 2013, the best-fit model poorly matches the data. Here, we investigate this year separately. The secondary spectra spanning 2013 are shown sequentially in Figure 5. Persistent clumps of power can be tracked throughout the year, and features can be seen to cross the line of f2 = 0, indicating a changing sign of vscr. From this, we can fit directly the signed value of 1/√θ = vscr (with the same free parameters s, Ψ, and vscr), where the locations of velocity cross-overs are quite constraining. The measures of vscr, and best fit model are shown in Figure 6.

The best fit screen parameters for 2013, and for all data beyond 2013 are tabulated in Table 1. The distance of the screen is consistent between both fits, with the orientation and parallel screen velocity differing between the two. This implies that the strong scattering plausibly arises from the same physical region. In addition, the absolute velocity of the screen need not be changing, as the orientation differs and we are only sensitive to the component of the screen velocity parallel to Ψ; the results of both fits are consistent with a screen velocity of |vscr| = 15 ± 2 km/s at an angle of φvel = 15° ± 10°.

We note however that the models above are incomplete, as a proper treatment needs to include the binary motion of the pulsar, which we have neglected. The orbit is 1.2 days, and vorb = 19.9 km s−1. Each observation is much smaller in duration than the orbit, but is at an effectively random orbital phase. This will add scatter in the velocities, and thus the curves. To properly account for the orbital velocity, one would need to jointly fit for i and Ω. Regardless, a 1D screen is a good fit to the data beyond 2013 (where each year the curvature peaks around November and is minimal around May), where the annual variation is the strongest effect.

### Table 1. Fit parameters of 1D screen to the arc curvatures, as defined in Section 5.2

|       | s (degrees) | vscr (km/s) |
|-------|-------------|-------------|
| 2014 onwards | 0.6 ± 0.06 | 106 ± 2 | −1.2 ± 2.5 |
| 2013 event | 0.58 ± 0.10 | 54 ± 9 | 12.8 ± 2.8 |

5.2 Location and nature of the scattering screen

As mentioned in 4.3, we fit the parabolic curvature of each observation to determine the masks for estimating the time delays. The arc curvatures contain the effective velocity, and vary throughout the year from the Earth’s motion. The existence of parabolic arcs suggests highly anisotropic scattering; for a one-dimensional screen,
5.3 Arclet evolution with Effelsberg: Scattering time on week to month timescales

We investigate the variability of scattering on week to month timescales with the Effelsberg observing campaign described in Section 3.2. The secondary spectra of all of our observations are shown in Figure 7. A clear parabolic arc, with a hint of inverted arclets is seen, and can be seen to clearly move through the secondary spectrum from left to right. We show three examples of larger, zoomed in secondary spectra in Figure 8 to emphasise these features. We estimate the total time delay from the secondary spectrum using the methods of Section 4.3, shown in the bottom panel of Figure 9. The total time delays are consistent with what was found...
with LEAP’s monitoring, showing steady scattering at around 60–100 ns, decreasing slightly over two months.

We also attempt to measure the effect of a single arclet, which is akin to contribution of a pulsar passing a single, compact point of scattering in the ISM, not unlike an echo. We track the position, and total fractional flux of the arclet seen travelling to the upper-right in the final 7 panels of Figure 7. This same feature first appears to be moving towards the origin in panels of MJD 58951-58958, although it is less prominent. We fit a flux centroid in an ellipse around the arc, to track its motion in $f_D = 0$ axis.

The motion of the arc unsurprisingly traces out a parabola over time, similar to echoes seen in the Crab pulsar (eg. Backer et al. 2000; Lyne et al. 2001). The strength of the arclet is quite asymmetric about the origin, and at its peak contains $\approx 4\%$ of the total pulsar flux. The total time shift arising from this arclet can be estimated as $\langle \tau \rangle_{\text{arclet}} \approx f_{\text{arclet}} / (f_{\text{arclet}}^2)$, and is shown in the middle panel of Figure 9. The contribution from a single arclet as it passes in front of the pulsar contributes a variable scattering of $\sim 20$ ns over a 2 month period.

5.4 Comparison to earlier results

Previous analysis using the ACF in Levin et al. (2016) measures the time delay from scintillation to be $\langle \tau \rangle = 11.7 \pm 4.9$ ns, monitoring this pulsar up until October 2013. Additionally, Shapiro-Albert et al. (2019) estimate a time delay of $\langle \tau \rangle = 43.6 \pm 2.3$ ns in a similar manner. The frequency channels used in these papers were 1.5625 MHz wide, a common standard in timing archives, and would have averaged over the fine scintillation structures due to power at high delays. Keith et al. (2013) estimate a scintillation bandwidth of 1.64 MHz, for which one would infer a time delay of 97 ns. This measurement is from before we have data, so we cannot directly compare, but this value is much closer to the order of the time delays we measure.

5.5 Effects of uncorrected scattering

In this section, we aim to make from our measurements a simple estimate of the single-pulsar gravitational wave signal arising from unaccounted scattering. We stress that this is a conservative upper-limit, as we do not know the extent to which variable scattering is absorbed in red-noise modelling, or in measurements of DM. The GW strain $h$ at a given periodicity $P$ is related to the amplitude TOA variations $\delta \tau$ as roughly $h \sim 2 \pi \delta \tau / P$. Since Aggarwal et al. (2019) find excess signal at 15 nHz, which is $\sim 2.1$ years, we estimate this using the long-term variable time delays from LEAP shown in Figure 4. We perform a Lomb-Scargle periodogram on the measured values of $\langle \tau \rangle$, and convert to a measure of $h$ while taking into account the proper normalizations, shown in Figure 10. The measured value at 15 nHz is $\sim 10^{-15}$, still an order of magnitude lower than the single pulsar limit of $h = 9.7 \times 10^{-15}$ from the EPTA (95% upper-limit, from Table 1 in Lentati et al. 2015). As PTA upper limits are improved, scattering variations, if unaccounted for, may begin to limit the timing precision.

6 CONCLUSIONS

We have measured variable time delays on a PTA millisecond pulsar, using similar methods to those laid out in Hemberger & Stinebring (2008). The next logical step will be to perform timing with scattering timescales subtracted from TOAs, to see if this improves the timing residuals. One can apply this approach to study the variable scattering in many PTA pulsars. With LEAP we can re-reduce the data to whatever time and frequency resolution we like, but regular timing observations would benefit with a second reduction with fine frequency channels at the expense of phase bins. Going further, methods to obtain the interstellar response directly may be important, including holography (Walker et al. 2008), cyclic spectroscopy (Demorest 2011; Walker et al. 2013; Palliyaguru et al. 2015; Dolch et al. 2020), or directly by using bright giant pulses in special cases (Main et al. 2017). Only these methods, in which the interstellar delays are measured directly (as opposed to measuring delay differences, as we do in the secondary spectrum) can retrieve overall delays that are not related to the characteristic timescale of the scattering tail. New analysis techniques such as the $\theta$-$\Theta$ diagram (Sprenger et al. 2020), which expresses the secondary spectrum in terms of the angular coordinates on the scattering screen, may be useful as well. This technique can be used to precisely measure arc curvatures, and could be used to efficiently perform holography of 1D screens (Baker et al. in prep.).

Scattering is statistically expected to follow $\sim \lambda^4$, yet scattering arising from discrete structures (observed as arclets, or localised clumps of power in secondary spectra) will be localised at a fixed $\tau$ as a function of wavelength. Observations over a wide frequency range will help to inform the amplitude scaling of arclets, and thus the contribution of discrete arclets to the total scattering time at different frequencies. In addition, scattering occurs from density gradients in the ISM, so the link between variable DM and scattering should be explored in more detail. In the case where the DM and scattering variations occur in the same scattering screen, they could potentially be inferred from the other quantity; a predictive model of scattering from DM (or vice versa) would be a great step towards removing these effects from timing observations.

Knowing the screen distance and orientation, the location of power in the secondary spectrum is predetermined, but the amplitudes are dependent on the physics of scattering and lens models. In recent years, some predictive models of scintillation properties have been developed (eg. Simard & Pen 2018; Gwinn & Sosenko 2019), which can be tested using measurements tracking arclets in time and frequency.
We used annual variations in the arc curvature to determine properties of the scattering screen, while orbital variations were ignored. Orbital variations can give an additional orbital constraint, such as the inclination (including the “sense”, Rickett et al. 2014; Reardon et al. 2019), and could possibly lead to precise pulsar distances (Boyle & Pen 2012). Such analysis would be greatly improved with precise, quantitative measurements of the arc curvature. Additionally, one could instantaneously measure the scattering screen’s properties using the multiple telescopes of LEAP, either using the visibilities (Brisken et al. 2010) or more simply using the inter-station time delays of the dynamic spectra (Simard et al. 2019a, while the method of combining of visibilities and intensities is outlined in Simard et al. 2019b). This is being investigated, and will be the subject of future work.

As scattering screens are likely much smaller than the angular separation between pulsars, scattering variations are very unlikely to directly correlate between pulsars in a way that mimics a Hellings & Downs curve (Hellings & Downs 1983). But while no direct correlation is expected between pulsars, it is possible that scattering is variable on similar timescales if pulsar proper motions and distances are comparable, and if the screen distance is not at an extreme (i.e. not too close to the pulsar or to the Earth). Several of the EPTA pulsars show variable scintillation arcs, similar to those shown in this paper, and will be subject of future work. As PTAs become more sensitive, any PTA result relying on a small number of pulsars may need to consider the effects of variable scattering when interpreting the significance of a potential gravitational wave signal.

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DATA AVAILABILITY

The timing data used this article shall be shared on reasonable request to the corresponding author.

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Figure 8. Zoomed in Effelsberg secondary spectra from Figure 7 for three different epochs, panels 1, 2 and 4. The red dotted line shows our best fit parabolic curvature, while blue dotted inverted parabolae are plotted with the same curvature, with their apex on the main parabola – these are to guide the eye towards faint structures in the secondary spectra which may be inverted arclets, and to help visualise the general trend of power moving from left to right along the parabola. More sensitive, or longer observations will be required to reveal the possible structure of inverted arclets more clearly.

Figure 9. Arclet evolution, and time delays from bi-weekly observations with Effelsberg. Top: Magnification (fraction of total intensity) and time delay of one arclet seen moving through the secondary spectrum. Red crosses are expected positions of the undetected arclet. Middle: contribution of the single arclet to the total time delay to the pulsar’s signal. Bottom: Estimate of the bulk scattering time inferred from the secondary spectra.

Figure 10. Estimated strain signal which would arise from uncorrected scattering, obtained from a Lomb-scargle periodogram of the bulk time delays (τ) from Figure 4.

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