Descattering of Giant Pulses in PSR B1957+20

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Abstract

The interstellar medium scatters radio waves, which causes pulsars to scintillate. For intrinsically short bursts of emission, the observed signal should be a direct measurement of the impulse response function. We show that this is indeed the case for giant pulses from PSR B1957+20: from baseband observations at 327 MHz, we demonstrate that the observed voltages of a bright pulse allow one to coherently descatter nearby pulses. We find that while the scattering timescale is 12.2 μs, the power in the descattered pulses is concentrated within a span almost two orders of magnitude shorter, of \(\lesssim 200\) ns. This sets an upper limit to the intrinsic duration of the giant pulses. We verify that the response inferred from the giant pulses is consistent with the scintillation pattern obtained by folding the regular pulsed emission, and that it decorrelates on the same timescale, of 84 s. In principle, with large sets of giant pulses, it should be possible to constrain the structure of the scattering screen much more directly than with other current techniques, such as holography on the dynamic spectrum and cyclic spectroscopy.

Key words: ISM: structure – pulsars: general – pulsars: individual (PSRB1957+20) – techniques: high angular resolution

1. Introduction and Background

When observed at relatively low frequency, pulsars scintillate, showing an interference pattern in frequency and time that arises from multi-path propagation through the interstellar medium. The interference patterns often have a clear structure, showing a quadratic dependence of the delay of scattering points as a function of their fringe rate (leading to parabolic arcs in the secondary spectra), which is most readily understood if the scattering is dominated by localized points on a strongly anisotropic screen (Stinebring et al. 2001; Walker et al. 2004; Cordes et al. 2006). This picture was confirmed dramatically by Brskens et al. (2010), who used very long baseline interferometry to show that the scattering screen of PSR B0834+06 appeared on the sky as a collection of points along a single, linear structure.

The above not only provides surprising information about the nature of the interstellar medium, but also offers a remarkable opportunity to study pulsars: given sufficient understanding of the locations of the scattering points, one can use them as an interferometer, which, with baselines of tens of au, has sub-microarcsecond angular resolution, thus allowing precision astrometry. Indeed, Pen et al. (2014) used the scintillation for PSR B0834+06 to show that the location of the radio emission shifted by a few tens of km as a function of spin phase.

The scintillation properties of pulsars are typically inferred from their dynamic spectrum, i.e., the intensity of a pulsar’s folded emission as a function of frequency and time. This yields direct information on the amplitudes of the interstellar impulse response function, but to retrieve its phase one has to rely on holographic techniques (Walker et al. 2008; Pen et al. 2014). A promising alternative is to use cyclic spectra of pulsars (Demorest 2011), which retain part of the phase information. So far, however, both methods have been shown to work only in specific cases.

In principle, the interstellar response could be measured directly if an object emitted bursts of emission that lasted much shorter than the scattering time. Some pulsars oblige by emitting suitably short “giant pulses.” One of these is the “black widow” pulsar PSR B1957+20 (Knight et al. 2006). In this paper, we show that its giant pulses indeed allow one to measure the interstellar response directly.

2. Giant Pulses

We recorded 9.5 hr of P-band data of PSR B1957+20 at the Arecibo Observatory, as part of a European VLBI network program (GP 052). The data were taken in four daily 2.4 hr sessions on 2014 June 13–16, recording dual circular polarizations of four contiguous 16 MHz wide bands spanning 311.25–375.25 MHz (recorded with the VLBA4 terminal in 2-bit Mark 4 format6). We exclude the fourth band from our analysis, as its signal was almost fully filtered out by the receiver, as well as the June 15 data, as these cover the eclipse of the pulsar by the wind of its companion (Fruchter et al. 1988), which hinders our analysis. We are thus left with 7.2 hr of data covering 311.25–359.25 MHz.

No flux calibrators were observed, so we convert to flux based on a nominal system temperature of 120 K and gain of 10 K/Jy for the 327 MHz receiver.7 With these values, the folded profile yields an average flux of 37 mJy, consistent with the 38 ± 3 mJy found by Fruchter & Goss (1992).

We searched for giant pulses in the de-dispersed timestreams of both polarizations, by binning the power in the whole 48 MHz band to 16 μs resolution and flagging peaks above 12σ (≈3 Jy). We use custom code to coherently de-disperse in 4 s blocks separately in each 16 MHz sub-band (using a dispersion measure of 29.1162 pc cm−3, tweaked using the folded profile),
and properly account for de-dispersion wraparound. We found 247 and 313 pulses in left and right-circular polarization, respectively, with 102 of these in common.8

In Figure 1, we show where the giant pulses arrive relative to the average profile. One sees three clusters, with the first (containing most pulses) coincident with the main pulse, the second coincident with the peak of the interpulse, and the third on the (relatively weak) tail of the interpulse. This distribution is different from what is seen in PSR B1937+21, where the giant pulses are predominantly found at the trailing edges of the pulse components (Cognard et al. 1996; Soglasnov et al. 2004).

The profiles of the giant pulses show a sharp rise and a long tail, suggesting that the pulses are intrinsically short. The tails are well fit by an exponential, with an e-folding timescale of 12.2 \( \mu s \) (see Figure 2, lower panel). The pulses are close to 100\% polarized, as expected for intrinsically short, single-mode emission. They show no preferred polarization direction, and can be strongly linearly or circularly polarized in either direction, in contrast to the folded profile, which is close to unpolarized (Fruchter et al. 1990).

3. Scintillation and Scattering

If giant pulses and the regular pulsar emission are both affected in the same way by propagation through the interstellar medium, an immediate expectation is that their frequency power spectra should show similar structure, and that this structure should vary on the same timescale. Comparing the power spectra for the brightest giant pulse with that of the regular emission near it, we indeed find that the spectra are very similar (see Figure 2). From the decorrelation bandwidth\(^9\) \( \Delta \nu = 133 \) kHz, we infer a scattering timescale \( \frac{1}{2\Delta \nu} \simeq 12.0 \mu s \), consistent with our measured timescale from the exponential scattering tail.

A stronger expectation is that the impulse response is the same, i.e., for giant pulses that happen close in time, not just the amplitudes of the impulse response function should match, but also the phases. We first verify this is the case for our closest pair, and then show it is possible to use a giant pulse as a direct measurement of the response function.

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8 At lower thresholds, many more true giant pulses are present, but separation from the bright tail of the “regular” pulses becomes more difficult, and these fainter pulses are less useful for our purposes here.

9 The half width at half maximum (HWHM) of the correlation function in frequency at zero time lag of the dynamic spectrum; e.g., Cordes et al. (1990).
Figure 3. Correlations between the voltage streams of a bright pulse pair separated by 1.92 s, for the three 16 MHz bands we observed in (with the two higher frequency bands offset by 1 and 2 units for clarity). This is equivalent to computing the visibility between the two pulses, as if they were a single pulse observed at two telescopes. The power is spread over ~200 ns, either because of small differences in the interstellar response, or, more likely, because of intrinsic differences in the two pulses. In either case, this spread sets an upper bound on the intrinsic duration of both pulses.

3.1. Giant Pulse Intrinsic Width

We directly compare the voltage timestreams of the two closest giant pulses, which are separated by 1.922 s. Both are strongly right-hand polarized, with an integrated signal-to-noise of 46 and 18, respectively. We compute visibilities between the two pulses as a function of time lag in each of the three 16 MHz, right-hand circular bands. In Figure 3, we show the visibility divided by the geometric mean of the autocorrelations, i.e., a measure of the correlation strength of the voltage streams. We find a strong correlation within a narrow envelope, which proves that the response of the interstellar medium is indeed close to identical for these two pulses, and that their intrinsic durations are very short, \( \lesssim 200 \) ns.

3.2. The Interstellar Response Function

The observed data are the convolution of the intrinsic electric field \( E_{\text{int}} \) with the impulse response function \( g \) of the interstellar medium, i.e., \( E_{\text{obs}}(t) = (E_{\text{int}} * g)(t) \). To the extent that giant pulses approximate impulsive intrinsic emission, it is thus possible to use them to measure the response, with which it should then be possible to undo the effects of scattering and scintillation, at least within the decorrelation time \( t_{\text{corr}} \), on which the response changes.

To verify this, we attempt to measure how well we can “descatter” pairs of giant pulses with each other. We do our analysis in Fourier space, where the pulses can be described as

\[
\hat{E}_{\text{obs}}(\nu) = \mathcal{F}(E_{\text{obs}}(t)) = \mathcal{F}((E_{\text{int}} * g)(t)) = \hat{E}_{\text{int}}(\nu) \hat{g}(\nu).
\]

(1)

If the intrinsic emission is a delta function at \( t = 0 \), one would have a Fourier spectrum with constant amplitude and zero phase, but because of the interstellar response, the signal gets mixed between frequency channels, causing amplitudes to change and phases to rotate (but with total power conserved).

In principle, the above suggests we could descatter a giant pulse by dividing by a suitably normalized reference pulse (as long as the two pulses occurred well within the timescale on which the interstellar response \( \hat{g}(\nu) \) changes, i.e., within \( \ll 84 \) s; see Section 3.3). In practice, this is rather noisy, as low-amplitude (and thus noisy) channels get upweighted, and high-amplitude (and thus well-measured) ones downweighted. An optimal solution to this would involve Wiener deconvolution, in which the different frequency channels are weighted properly using their signal-to-noise ratio. We opted instead, however, to just normalize the reference pulse by its amplitude, i.e., use only the phase information. Beyond simplicity and the knowledge that most of the information is contained in the phases (in terms of power, a fraction \( \sim \pi/4 \); see Section 3.3), this has the advantage of being complementary to the dynamic spectrum, which considers only the amplitudes. Specifically, dividing a trial pulse by a normalized reference pulse gives

\[
\frac{\hat{E}_{\text{trial}}}{\hat{E}_{\text{obs}}(\nu)} = \frac{\hat{E}_{\text{trial}}}{\hat{E}_{\text{int}}(\nu)/\hat{g}(\nu)} \frac{\hat{g}(\nu)}{\hat{g}(\nu)}
\]

(2)

where we dropped the dependence on \( \nu \) for brevity.

If the pulses have the same response function and the reference pulse is truly impulsive, this reduces to \( \hat{E}_{\text{int}}(\nu) / \hat{g}(\nu) \), i.e., one would recover the intrinsic spectrum of the trial pulse multiplied by the amplitudes of the response function. Since the phases are corrected but the amplitudes are not, for an intrinsically short pulse, an inverse Fourier transform should yield a timestream with a pulse that is similarly short but has reduced amplitude. More generally, since we do not have perfect arrival times, there will be an uncertainty in the time offset (equivalent to a phase gradient in the spectrum), and, since the intrinsic pulses are not true delta functions, the power will only be concentrated within the intrinsic width of the emission.

We apply our method first using our brightest giant pulse as a reference, and its five closest neighbors as trial pulses. Our brightest pulse is strongest in right-circular polarization, with an integrated signal-to-noise of 132. We make Fourier transforms for 32 \( \mu \)s segments for all pulses, covering the majority of the scattering tail (this corresponds to 1024 real-valued samples and thus 512 channels in each of the three 16 MHz bands). We then descatter and inverse transform as above, and bin and sum the power of the descattered timestreams in 250 ns bins to account for possible intrinsic widths (see Figure 3).

The results are shown in Figure 4. For the closest pair, one sees that much of the power of the descattered pulse is contained within a single 250 ns bin, and that the peak intensity of the descattered pulse is more than 10 times stronger than that of the observed, scatter-broadened one. For the next closest pulse, the procedure does not seem to work well, with the descattered pulse having multiple peaks. This is likely intrinsic, since for the further pulses, the descattering does work, though with decreasing efficiency, presumably because the response function becomes increasingly different.

3.3. Decorrelation Time

The above suggests that it should be possible to use the extent to which giant pulses can descatter each other to determine the timescale on which the response function changes. For this purpose, we increase our sample by selecting as reference pulses all those for which phases can be measured in the Fourier spectrum, i.e., with \( \gtrsim 1 \sigma \) per voltage sample, corresponding to an integrated signal-to-noise of 40 in the power. This results in six and three pulses in the right and left...
circular timestreams, respectively (with one in common). For each of these, we descatter all pulses within five minutes to either side, matching the polarizations. We then sum the power of the descattered timestreams in 1 µs bins (which should account for any reasonable intrinsic widths), and measure the fraction of the total power which is descattered into a single peak to quantify how well the pulses descattered each other.

With our brightest pulse, we find that we can recover up to 70% of the total flux in the descattered peak. For the other, fainter pulses, however, at most half of the power is descattered, as their more limited strength allows only an imperfect measure of the response function. To get a quantitative sense of the imperfections, we test our routine on simulated giant pulses. For the simulations, we use the following two assumptions (which are clearly simplistic but should cover the essence of short intrinsic emission and an exponential scattering tail): that at our time sampling the intrinsic emission of giant pulses can be represented by delta functions, and that the response function \( g \) can be described by a normally distributed random process for which the variance decreases exponentially on the observed 12.2 µs decay time (normalized to have unity integrated power). Hence, our sets of simulated giant-pulse timestreams consists of sets of delta functions of different amplitudes convolved with a given simulated response (of length 32 µs, i.e., 3072 real-valued samples), with normally distributed measurement noise added to each sample. We run these pulse pairs through our analysis routines to determine the fraction of power that can be descattered, given two pulses with identical response functions.

As expected, the strength of the reference pulse determines the average power recovered, while the strength of the (fainter) trial pulse dominates the scatter between different realizations. For our set of reference pulses, we find that the simulated recovered fraction ranges from 0.4 to 0.7. From Equation (2), one sees that in the high signal-to-noise limit the described fraction should be dominated by the extent to which \( \langle |g| \rangle^2 \) is less than one. Since our simulations assume \( 2g^2 \) is distributed as a \( \chi^2 \) distribution with 2 degrees of freedom, one expects \( \langle |g| \rangle^2 \approx \pi/4 \), consistent with our results. For lower signal-to-noise, we find that we can also reproduce our simulated fractions by integrating numerically over probability distributions for both \( \bar{g} \) and the noise (unfortunately, we could not find a closed-form expression).

In Figure 5, we show the fraction of the power that is descattered against time separation for each pulse pair, corrected for the above loss of power (and with error bars reflecting the expected 1σ scatter). One sees that at large separation, \( >100 \) s, the descattering never recovers much power, while at shorter separation it does, though unequally so. Inspection of the low points around \( \Delta t \approx 20 \) s shows that in those cases the descattering does not lead to a single strong pulse (as for the \( \Delta t = 35.8 \) s pulse in Figure 4). This likely reflects intrinsic pulse structure, in either the descattered pulses or the reference pulse. Three of the nine reference pulses descatter a pulse pair to within a single 1 µs peak (including our brightest two pulses, shown in Figures 3 and 4), four of our reference pulses show clear success in descattering adjacent pulses but with remaining multi-peaked structure, and the last two reference pulses have no pulse pairs within 84 s separation. We do not have a sufficient number of pulse pairs within the decorrelation timescale to determine which, if any, of our reference pulses are intrinsically wide.

We can compare our decorrelation timescale with that derived using the more traditional way, from the autocorrelation of the

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**Figure 4.** Descattering giant pulses. Top: the profile of the brightest giant pulse in our sample, used to descatter its neighboring pulses. Lower panels: nearby giant pulses, ordered by time separation, showing both the profiles of the observed, scatter-broadened pulse (red curve, lowered by 50 Jy), and the descattered pulse (black curve). For all pulses, the signal is binned to 250 ns. One sees that the descattering works less well at larger time differences, for which the interstellar response starts to change significantly. The outlier at \( \Delta t = 35.8 \) s appears to be formed of a series of bursts, in contrast to the other pulses, which appear to be well described by delta functions in this time binning.
of the dynamic spectrum, the correlation coefficient is
\[ R(\tau) = \frac{\langle (I(t, \nu) - \mu)(I(t + \tau, \nu) - \mu) \rangle}{\sigma^2}, \]
where \( \mu = \langle I(t, \nu) \rangle \), \( \sigma^2 = \langle (I(t, \nu) - \mu)^2 \rangle \), and \( \langle \rangle \) denotes an average over both time and frequency.

We apply this to a dynamic spectrum created for the nine minutes of data surrounding our brightest giant pulse. We use bins of 4 s, 8 kHz, and 1/32 in phase (where the frequency and phase resolution are set to barely resolve the main pulse; at 8 kHz, we also barely resolve the frequency structure due to scintillation). We define two off gates, with one subtracted from the folded profile to give a pulsed flux, and the other used as an independent measure of the noise in the dynamic spectrum. We then compute the autocorrelation of the dynamic spectrum, subtracting the autocorrelation of the noise, and averaging over the central 14 MHz of each band (to avoid the parts most affected by bandpass variations).

The result is shown in Figure 5. One sees that at short times the correlation is very good (it approaches 0.98 rather than 1, likely because we ignore the frequency dependence of the noise), and then it decreases smoothly. Taking the decorrelation time as the lag where the correlation drops by 1/e, we find \( t_{\text{corr}} = 84 \) s.

In Section 3, we showed that the spectrum of brightest pulse was similar to that of the regular emission. With the dynamic spectrum, we can verify this quantitatively, and also check the dependence on lag. For this purpose, we Fourier transform our brightest giant pulse to the same channelization as the dynamic spectrum (where 8 kHz corresponds to 125 \( \mu \)s in time, i.e., it covers the full width of the scattering tail; see Figure 2). We then correlate it against the dynamic spectrum for a range of delays.

At low delay, we find that the giant pulse correlates very strongly with the dynamic spectrum, at 92 \pm 5\% (see Figure 5). This strong correlation independently suggests a short intrinsic duration of the giant pulse, as intrinsic structure on timescales comparable to the scattering time will lead to differences in the spectra (as seen for the scintillation pattern of the Crab’s giant pulses, Cordes et al. 2004). At longer delays, the correlation drops, and we find \( t_{\text{corr}} = 84 \) s, the same value obtained from the autocorrelation of the dynamic spectrum.

Comparing the points from descattering giant pulses with each other, one finds that the curves derived using the power spectra form a rough upper envelope, suggesting that the decorrelation time is the same. This is not a trivial comparison, since the autocorrelations take into account the amplitudes of the impulse response function only, while our descattering only uses the phase information.

4. Ramifications

We have shown that giant pulses in PSR B1957+20 can be used as a direct probe of the impulse response function of the interstellar medium: they allow one to descatter other giant pulses. An immediate result is that we can constrain the typical intrinsic duration of giant pulses of PSR B1957+20 to be very short, \( < 200 \) ns.

Having a direct measure of the response also should allow one to verify (and inform) the response inferred from the dynamic spectrum (using holographic techniques; Walker et al. 2008; Brisken et al. 2010) or from cyclic spectra (Demorest 2011; Walker et al. 2013), thus putting those techniques on firmer footing.

One might also hope to use the giant pulses directly to infer how the impulse response changes with time. In general, the decorrelation time is only an average estimate; parts of the impulse response at short delay should change more slowly than parts at large delay. Indeed, in the longer term, one might hope to use all giant pulses—perhaps aided by the dynamic and cyclic spectra—to measure the evolution of the response as a function of time and frequency. If the scattering toward PSR B1957+20 is dominated by a single, highly anisotropic screen, as was found for PSR B0834+06 (Brisken et al. 2010), then this should allow one to determine amplitudes and phases of individual scattering points directly. Those, in turn, might allow one to resolve the pulsar’s orbit on the sky, as has been done for the pulse emission with spin phase for PSR B0834+06 (Pen et al. 2014).

Unfortunately, only a few other pulsars show giant pulses. Among those, we are most excited to apply our technique to the Crab pulsar, since for that source we expect that the main scattering screen, which is in the Crab Nebula, will resolve the light cylinder, thus opening up the possibility to determine empirically where the giant pulses originate.

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Facility: EVN:Arecibo:327-MHz Gregorian.

Software: Astropy (Astropy Collaboration et al. 2013).

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