Efficient light storage in a crystal using an atomic frequency comb

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Abstract. We demonstrate efficient and reversible mapping of a light field on to a thulium-doped crystal using an atomic frequency comb (AFC). Owing to an accurate spectral preparation of the sample, we reach an efficiency of nine per cent. Our interpretation of the data is based on an original spectral analysis of the AFC. By independently measuring the absorption spectrum, we show that the efficiency is limited by both the available optical thickness and the preparation procedure at large absorption depth for a given bandwidth. The experiment is repeated with less than one photon per pulse and single-photon counting detectors. We clearly observe that the AFC protocol is compatible with the noise level required for weak quantum field storage.

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1. Introduction

The ultimate control of a quantum light field naturally demands a good understanding of its interaction with matter. In that context the reversible mapping of a field into an atomic system represents a fundamental challenge. The experimental quest for a quantum memory for light has produced many achievements using different systems. This domain has also been motivated by the realization of a quantum repeater for which memory is a critical element [1]. Benefitting from long experience, most of the recent achievements have involved atomic vapors [2]–[6]. More recently, atom-like impurities in solid-state matrices have been considered to be excellent candidates for a quantum interface. They can indeed have good coherence properties [7, 8]. One can take advantage of their complete absence of motion to demonstrate extremely long coherence time [9]–[11]. The rare-earth ion doped crystals (REIC) are particularly interesting because they are relatively easy to produce and are commercially widely available. The most promising storage protocols fully use the primary feature of REIC: the optical transition has a long coherence lifetime and a large inhomogeneous broadening. These schemes are in the lineage of the photon echo technique and are intimately related to a dipole rephasing. The direct excitation of the optical coherences over a wide spectral range (larger than the homogeneous width) is usually the initial storage step. The role of the protocol is to rephase the dipoles by externally controlling the detuning of the transition [12]–[15], which can lead to large storage efficiencies [16].

More recently, a protocol based on the preparation of an atomic frequency comb (AFC) has allowed light storage at the single-photon level [17, 18]. The protocol is based on the filtering of a probe photon by a spectral array of evenly spaced narrow absorption peaks. The resulting absorption comb of periodicity $1/T$ is prepared in advance through the shaping of the population spectral distribution. The weak signal pulse, much shorter than $T$, is sent into the medium. Its spectrum covers many comb peaks. The comb gives rise to multiple delayed responses, evenly spaced by the interval $T$, as expected from the frequency-to-time conjugation properties of periodic spectral filtering. In a similar way, a spatially periodic grating diffracts an incident beam in evenly spaced directions, because angle and spatial coordinate are conjugated variables. Within the frame of this analogy, the AFC response at delay $T$ plays the role of the first diffracted order. This simple analogy does not reveal the complexity of the original protocol.
of AFC storage, even if it helps in understanding the retrieval mechanism at time $T$. The AFC response also includes causality, which forbids the emission of negative orders and, in association with propagation features, determines the relative weight of successive orders at times $T$, $2T$, $3T$, . . . . We only consider the first order (AFC echo at time $T$). We shall take the causality into account to calculate the efficiency of the retrieval in the AFC echo (section 3.1).

This spectral image has been widely used to explain the stimulated photon echo (SPE). The AFC is in the direct lineage of the SPE where a periodic modulation of the absorption profile is considered. In that sense, the AFC represents an extreme situation: high narrow absorbing peaks are separated by a fully transparent region. Such a comb maximizes the retrieval of the incoming energy into the AFC echo at time $T$. As demonstrated by Afzelius et al [17], the retrieval efficiency should be high in the forward direction (54%) and perfect (100%) in the backward configuration. In order to achieve high efficiency, it is crucial to achieve a precise spectral shaping of the inhomogeneous absorption profile. The AFC is not only an interesting alternative storage protocol, it has also been shown theoretically to store the highest number of pulses for a given bandwidth [19]. It defines the number of temporal modes and is critical for the prospect of quantum communication [20].

In this paper, we show how the AFC protocol can be efficiently implemented in an appropriate crystal, namely Tm$^{3+}$:YAG. Our work is focused on the initial step of the protocol where the signal is stored on to the optical coherences. Our results in terms of storage efficiency correspond to more than an order of magnitude of improvement on the previous proof-of-principle realization performed in Nd$^{3+}$:YVO$_4$ [18]. This improvement is possible because Tm$^{3+}$:YAG allows for a better preparation of the AFC due to efficient optical pumping and narrower homogeneous linewidth. More generally, the material properties give the final width of the spectral selection. It can be limited ultimately by the interaction with neighboring ions (from the crystalline matrix or rare-earth dopant). Additionally, any external source of broadening such as the laser linewidth can reduce the efficiency for a given bandwidth. This limiting factor is not present in our case because we use a stabilized laser.

2. Efficient light storage in Tm$^{3+}$ : YAG

Thulium has an interaction wavelength (793 nm) more easily accessible with laser diodes than other rare-earth ions (praseodymium and europium). This is particularly convenient when a complex spectral preparation sequence is required. The trivalent thulium is also known to have a long optical coherence time from 10 to 100 $\mu$s, depending on the crystalline matrix, dopant concentration and magnetic field [21]. The spin coherence lifetime has also been observed to be long in the ground state [11]. Very precise spectral tailoring will ultimately lead to a high number of temporal modes stored in the medium, which primarily scales as the number of spectrally independent peaks in the absorption profile (ultimately limited by the homogeneous width). Optical pumping to a shelving state offers an effective way to shape the population distribution of the absorbing atoms [22]. In Tm$^{3+}$:YAG, the nuclear spin states in the ground state can be split by an external magnetic field and one of them can be used as the optical pumping shelving state. Those split states exhibit a long lifetime that makes them adequate for a complex initial state preparation [23, 24].
2.1. Experimental set-up

In Tm\(^{3+}\):YAG the electronic state degeneracy is completely lifted by the crystal field. Therefore, unlike free atoms in a vapor, the impurity ions hold optical dipoles with a precisely defined spatial orientation. The narrow line transition dipole moments connect the singlet states of lower energy in Stark multiplicities. With a nuclear spin \(I = 1/2\), Tm\(^{3+}\) does not exhibit hyperfine splitting, unlike other rare-earth ions such as Eu\(^{3+}\) or Pr\(^{3+}\). The hyperfine interaction is only revealed to second order, through coupling with the electronic Zeeman Hamiltonian. In other words, the nuclear Zeeman splitting is enhanced by second-order hyperfine coupling with the electronic Zeeman effect. The coupling is different in the two electronic states connected by the optical transition. As a consequence, the nuclear spin can be flipped by an optical excitation, in contradiction with the usual nuclear spin selection rule \([25]\).

To implement the AFC protocol in Tm\(^{3+}\):YAG, we apply a 210 G magnetic field along the [001] crystalline axis. This splits the ground and excited levels into a nuclear spin doublet \((m_I = \pm 1/2)\) by \(\Delta_g = 6\) MHz and \(\Delta_e = 1.3\) MHz, respectively (see figure 1). These values are relatively standard for rare-earth ions with an even number of electrons in a low-symmetry site \([26]\). The ground state splitting is sufficient to create an interaction bandwidth of a few megahertz (\(\sim 3\) MHz in our case).

Our 0.5% doped Tm\(^{3+}\):YAG crystal is immersed in liquid helium at 2.3 K. This permits a good coherence lifetime \((T_2 = 30\) µs) and then an accurate spectral tailoring. This is also required to tailor the spectral distribution of populations by shelving them from one nuclear state to the other (comb preparation). We measure the population shelving time to be 7 s in our case. We also observe that it depends critically on the temperature \([23]\). Although small (typically 2% for our field orientation \([25]\)), the optical branching ratio is compensated for by the long
shelving time. This is sufficient to efficiently pump the ions between the nuclear spin levels. The spin mixing can also be assisted by the non-radiative decay through the \(^3\text{F}_{4}(0)\) bottleneck state (10 ms population lifetime). An evaluation of this effect requires further investigation of the optical pumping dynamics and cannot be deduced from our measurements. The laser polarization is also applied along the [001] axis. In this configuration, the splitting is the same for the different crystallographic sites (orientations of the transition dipole \([25]\)). In the present set-up the magnetic field is not strong enough to generate a large \(\Delta_{e}\) splitting, which would push the shelving state outside the AFC operating bandwidth. This can be an issue, in practice, because burning a periodic ion distribution gives birth to two replicas shifted by \(\pm \Delta_{e}\). In order to avoid interference between these multiple AFCs, we simply make the comb period \(1/T\) coincide with the excited level splitting \(\Delta_{e} = 2/T\) in our case).

The spectral preparation is performed by applying a sequence of pulse pairs (full-width at half-maximum (FWHM) duration \(\sim 300\) ns separated by \(T = 1.5\) \(\mu\)s and then followed by a 100 \(\mu\)s dead time before the next pair). This train consists of 5000 pairs followed by a long waiting time (50 ms) to ensure the complete decay of population from the \(^3\text{H}_{4}(0)\) excited state (lifetime \(\sim 800\) \(\mu\)s). Each pair whose area is much smaller than \(\pi\) creates a small population difference proportional to the optical power spectrum at this specific detuning. In this limit, the preparation sequence can be seen as frequency-selective optical pumping with a \(1/T\)-periodic optical power spectrum. This interpretation is valid as soon as the separation between successive pairs (100 \(\mu\)s) is longer than the optical coherence time. Accumulation by repeated illumination is the very essence of optical pumping. This is the only way to circumvent the efficient decay of the upper level to the initial ground state and to ultimately favor the spin-flipping weak transition. After accumulation of many pairs, the sequence gives rise to the AFC. An alternative interpretation of the pulse pair has been proposed by Afzelius \textit{et al} \([18]\). A complete study of the atomic dynamics during the preparation stage is beyond the scope of the present paper (since we focus on the storage of the resulting population structure) but deserves further investigation \([27]\).

2.2. Atomic frequency comb spectrum and storage efficiency

To first probe the population frequency spectrum and then generate a pulse to be stored, we use a very weak independent signal beam. Temporal shaping and frequency scanning are achieved by two independent acousto-optic modulators (AOM). The pumping beam has also been designed to be twice as large as the signal in the crystal. The AFC is then spatially uniform over the entire signal beam.

2.2.1. Probe transmission spectrum. By sweeping the signal AOM driving frequency, we can directly observe the transmission spectrum (see figures 2(a) and (c)). Even if the pumping power is very weak, we see that the AFC structure depends strongly on the preparation pulse intensity. For an increasing preparation power from figure 2(a) to figure 2(c), the peak maximum absorption decreases due to power broadening of the hole-burning process (see insets). The sweep rate \((320\) kHz\(^2\)) being slightly too high, we observe a distortion of the observed spectrum (see figure 2(c) for example). This effect is not apparent in the insets representing only the central part of the spectrum where a slower sweep rate \((120\) kHz\(^2\)) has been used.

2.2.2. Weak signal storage. A weak pulse to be stored is then sent on to the medium. To be sure that the signal is uniformly covered by the comb, its duration is chosen to be slightly
Figure 2. (a) and (c) Probe transmission spectrum obtained for different preparation pulse powers. The 660 kHz period is given by the delay between the preparation pulses $T$. This structure is tailored within the inhomogeneous profile ($\approx 10$ GHz). The transmission goes to zero out-of-resonance, which is consistent with the large initial optical depth of our sample ($\approx 4–5$). The insets show the central part of the optical depth spectrum. (b) and (d) A weak probe signal is sent matching the AFC bandwidth of (a) and (c), respectively (this reference is called ‘incoming’ in the figure). After being partially absorbed (called ‘transmitted’), it gives rise to a retrieval (called ‘retrieved’). The efficiency strongly depends on the shape of the comb, which in turn depends on the power of the preparation pulses. The peak preparation power is weak (fraction of $\mu$W) and is not monitored independently in our set-up. The preparation power is larger by one order of magnitude typically for the set (c)–(d) than for (a)–(b).
longer than the preparation pulses (FWHM $\sim 450$ ns). The transmitted signal gives rise to a delayed retrieval. Curve 2(b) (resp. 2(d)) corresponds to 2(a) (resp. 2(c)). To be able to compare the incoming pulse and the retrieved signal, we apply a strong single pumping beam (a thousand times stronger than the preparation pulses). It burns a broad hole (10 MHz width, typically). Since all the population has been pumped out of the hole center, the probe is considered as transparent (photo-bleaching). This procedure gives a reference for a fully transmitted incoming pulse. We can then estimate the total efficiency: $\eta_{\text{exp}} = 9.13 \pm 0.10\%$ (resp. $\eta_{\text{exp}} = 4.37 \pm 0.09\%$) for the comb 2(b) (resp. for 2(d)). This corresponds to more than one order of magnitude of improvement on the previous proof-of-principal realization [18].

3. Discussion

Because of the independently measured AFC spectrum, we now investigate storage efficiency. To do so, we have developed an original model to analyze our data. The expected efficiency has already been calculated by assuming that the comb is composed of well-separated Gaussian peaks [17]. This calculation was based on a temporal image by following the evolution of the atomic variables during the time sequence. Here we derive a more general formula valid for any peak shape.

3.1. Expected efficiency

We consider the propagation of a field $\mathcal{E}$ in the spectral domain:

$$\partial_z^2 \mathcal{E}(z, \omega) + k^2 (1 + \chi(\omega)) \mathcal{E}(z, \omega) = 0.$$ 

The action of the medium is represented by the susceptibility $\chi(\omega)$ ($k$ is the wave vector). To account for the periodic structure of the AFC (period $1/T$), we naturally decompose $\chi(\omega)$ in a Fourier series:

$$\chi(\omega) = \sum_{p \geq 0} c_p \exp(-2i\pi p \omega T).$$

The components corresponding to $p < 0$ are dropped for causality reasons. We here calculate the efficiency in the forward direction, but our analysis could be extended to describe the backward retrieval [17]. The field $\mathcal{E}(z, \omega)$ has the same periodic structure:

$$\mathcal{E}(z, \omega) = \mathcal{E}(0, \omega) \sum_{p \geq 0} a_p(z) \exp(-2i\pi p \omega T).$$

Each term of the decomposition gives the amplitude of a retrieval centered at the time $pT$. The first term $p = 0$ represents the absorption of the incoming signal. In the forward direction, the efficiency of the AFC protocol is then directly given by $|a_1|^2$. All the coefficients can be recursively determined by solving the propagation equation for $a_p(z)$. The first two components are given by

$$\partial_z^2 a_0(z) + k^2 (1 + c_0) a_0(z) = 0,$$

$$\partial_z^2 a_1(z) + k^2 (1 + c_0) a_1(z) + k^2 c_1 a_0(z) = 0.$$
The efficiency in the forward direction is then given by
\[
\eta = \frac{1}{4} \frac{|c_1|^2}{\text{Im}(c_0)^2} \tilde{d}^2 \exp(-\tilde{d}).
\]
(1)

\(\tilde{d}\) is the average absorption depth defined as \(\tilde{d} = -k \text{Im}(c_0) L\), with \(L\) being the length of the medium. Our formula is very general and simply assumes that the susceptibility is periodic. It is sufficient to estimate the expected efficiency for a given frequency comb as presented in figure 2. The main advantage of our model is the absence of assumption on the exact shape of the profile. In addition, one gains an alternative interpretation in the spectral domain; the periodic filter is considered as a diffraction grating. A simple numerical Fourier expansion of the absorption spectrum (see insets of figures 2(a) and (c)) gives the different coefficients to determine \(\eta\). We calculate \(\eta = 7.8\%\) (resp. \(\eta = 3.5\%)\) for figure 2(a) (resp. 2(c)). More generally for different preparation pulse intensities, we find a correct agreement between the experimental efficiencies and those deduced from the measured absorption spectrum (\(\eta_{\text{exp}}\) is represented by a red dashed line and \(\eta\) by a black dashed line in figure 3). Nevertheless, we systematically observe that the measured values are larger than predicted. This may be due to inaccurate measurement of the absorption spectrum. As already mentioned, the lower sweep rate 14.4 kHz \(\mu s\) used to obtain the central part of the spectrum may still be too high compared to the width of the AFC peaks. This acts as a convolution effect, artificially broadening the peaks. More generally, it is not fully obvious to accurately measure an optical thickness much larger than one. Any very distant resonance contribution of the laser would be transmitted through the medium. Here, it could

**Figure 3.** Measured efficiency \(\eta_{\text{exp}}\) as a function of the maximum optical depth (red dashed line used to guide the eye). We have also calculated the expected efficiency \(\eta\) (equation (1)) from the absorption spectrum (black dashed line used to guide the eye). Then assuming that the comb is composed of Lorentzian peaks (equation (2)), we have finally plotted the optimum efficiency \(\eta_{\text{opt}}\) (solid black line) and fitted the experimental spectrum to obtain the finesse. Inset: measured finesse obtained by fitting the absorption spectrum with a comb of Lorentzian peaks (red dashed line used to guide the eye) and comparison to the optimum \(F_{\text{opt}}\) (solid black line). The two cases presented in figure 2 are circled.

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be due to the additional tapered amplifier whose fluorescent background is extremely broad. For both effects, the real optical thickness would be underestimated, as would the predicted efficiency.

3.2. Current limitations of the pumping scheme

It is now important to understand the current limitations. We first estimate the maximum efficiency that could be obtained in our case and compare it to our measurements (see figure 3). The total efficiency strongly depends on the exact shape of the frequency comb. For a given initial optical thickness, one has to properly choose the peak width to optimize the retrieval. We now assume that the comb is composed of Lorentzian (half-width at half-maximum $\Gamma$) periodically distributed peaks. The Lorentzian shape allows simple analytical calculation and gives a satisfying agreement with the measured spectrum. The efficiency depends only on the maximum optical depth $d$ and the comb finesse $F = \pi/(\Gamma T)$ [17]:

$$\eta = d^2 \tanh^2(\pi/2F) e^{-d \tanh(\pi/2F)} e^{-2\pi/F}. \quad (2)$$

This result can be compared with equation (1) of [18]. The structure and the term interpretation are similar. Our formula includes, by definition, a potential overlap of the Lorentzian peaks when the finesse is low. This has been alternatively interpreted by de Riedmatten et al [18] as a uniform absorbing background and fitted independently. In our case, we can directly fit the AFC spectrum with equation (2). Our result can also be simplified when the finesse is high. In that case, the best efficiency of the protocol is expected. For a given initial optical thickness $d$, the finesse should be properly chosen to obtain an optimized efficiency. For $F_{\text{opt}} = \pi(1 + d/4)$, we have

$$\eta_{\text{opt}} = 4e^{-2d^2/(4 + d)^2}. \quad (3)$$

The efficiency is limited to 54% in the forward direction. This value is a common feature of the protocols involving a dipole rephasing [15].

One can now consider whether the measurements correspond to the optimum for a given optical depth. As we have seen, the shape of the comb is essentially controlled by varying the intensity of the preparation pulses. At the same time this procedure is changing the maximum absorption depth $d$ and the width of the peaks (the finesse $F$). For different pumping power, we can directly measure the maximum absorption depth $d$ and then estimate the corresponding $\eta_{\text{opt}}$ (equation (3)).

The comparison between the measured efficiencies $\eta_{\text{exp}}$ (red dashed line in figure 3) and the expected optimum for a given optical depth $\eta_{\text{opt}}$ (solid black line in figure 3) allows us to evaluate the quality of the preparation. There is clearly a good matching of the efficiency curves at low optical depth (between 0 and 3). In that case, the efficiency does not critically depend on the finesse (equation (2)) [17]. In practice, any finesse between $\pi$ and $2\pi$ could be considered as optimal. This is verified experimentally and explains the efficiency curve matching even if the finesse only approximately follows $F_{\text{opt}}$ (see inset in figure 3). On the contrary, we observe an efficiency breakdown at larger optical thickness ($d \sim 3$ in figure 3), where $\eta$ now critically depend on the finesse (equation (2)). The measured finesse indeed significantly deviates from $F_{\text{opt}}$ (see inset in figure 3). This is sufficient to explain the efficiency breakdown in this region.

This clearly shows that the preparation sequence may not be appropriate at large optical thickness. Even if we are able to benefit from the initial optical depth of our crystal ($d \sim 4 - 5$), we cannot properly prepare the AFC to obtain the maximum efficiency. In practice, the external
control of the preparation is given by the power of the pumping pulses. This single degree of freedom cannot independently master the two parameters of the comb (height $d$ and finesse $F$). Our comparison shows that a simple pumping strategy using a train of pulse pairs is most appropriate at low optical thickness as soon as the power is properly adjusted. By contrast, at large optical thickness the efficiency is limited.

The pumping intensity primarily affects the width of the hole-burning process because of power broadening. This effect is intimately related to optical pumping dynamics. To break this limitation, it is possible to apply a coherent population transfer scheme [28]–[30]. This approach has been successfully implemented to obtain and isolate narrow features in different systems, including Tm$^{3+}$:YAG [29]–[31]. In other words, by enriching the preparation sequence with numerous pulses, one would obtain another degree of freedom to accurately control the comb shape. The propagation of sophisticated pulse sequences in highly absorbent materials may nevertheless limit the capability of the coherent population transfer. In that case, one can also use a transverse illumination of the slab [32]. Combining these two techniques, it should be possible to independently control the efficiency and spectral properties of the AFC.

4. Storage at the single-photon level

Even if efficiency is a critical parameter for a quantum storage protocol, a noise estimation should also be performed. In order to investigate this level in the very weak field condition, we reproduce the experiment with less than one photon per pulse.

We use the same preparation procedure that has been optimized previously. The weak incoming pulse to be stored is strongly attenuated by using neutral density filters. After transmission through the crystal, the signal is now fiber coupled to a single-photon counting module (Perkin-Elmer SPCM AQR 14 FC). To increase the measurement statistics, we repeat 1000 signal pulses after each preparation procedure (5000 pairs followed by a 50 ms waiting time). The total duration of the sequence is now 329 ms corresponding to 3039 signal pulses per second. The arrival of the photon is counted in a 300 ns time gate and accumulated for 5.51 s. The corresponding histogram is presented in figure 4 for different positions of the time gate covering the arrival of the transmitted signal and the AFC echo.

We have independently measured the collection efficiency from the crystal to the single-photon counting modules. Knowing the quantum efficiency of the detector and assuming that the AFC efficiency is the same as previously measured, we estimate the incoming number of photons per pulse in the crystal to be $\sim 0.5$.

For comparison, the measurement has also been carried out at higher power (solid curve in figure 4). We observe the same kind of trace as previously described (see figure 2 for example). Nevertheless, we suspect the efficiency to be higher because we also observe a second echo. This may be due to the slightly different experimental conditions (magnetic field). We have not measured the efficiency in that case, but this effect is a potential route to improve retrieval and is currently under investigation. The photon counts (histogram in figure 4) follow the same curve within the statistical fluctuations. We additionally observe a detection background of $\sim 8$ counts. By comparing this noise level to the maximum of the retrieved signal, we obtain a signal-to-noise ratio of 11. The intrinsic dark counts of the detector should be 0.31 with this cycling, gating rate and accumulation time. We have isolated this detection background, which is due to a leak of the preparation beam. Even if we use a first AOM to temporally shape the preparation pulse and a second one right before the detector, the total isolation is not sufficient to completely...
Figure 4. Histogram of the accumulated counts for 5.51 s within a 300 ns time gate. The incoming pulse contains \( \sim 0.5 \) photons per pulse. We have recorded independently the signal at higher power using a standard photon diode (solid red curve). Three observed pulses are, respectively, the transmitted power, the retrieved signal (first echo) and a second echo (see text for details).

remove the residual light. In the near future, we can also use a mechanical chopper since the waiting period is long (50 ms). Finally, we have not seen any spurious light due to the medium itself (fluorescence). The spontaneous emission is strongly reduced after a long waiting time and is also emitted in all directions. We therefore believe that the ultimate limitation of such a measurement involves the dark counts or electrical afterpulses of the single-photon counters. In other words, the AFC storage should not add any noise to the incoming photons. This is a remarkable advantage that makes the protocol promising for the storage of a non-classical field.

5. Conclusion

To conclude, we have successfully applied the AFC protocol to a Tm\(^{3+}\):YAG crystal whose properties are of particular interest for the prospect of quantum repeaters. We have observed efficiencies up to 9.13 \( \pm \) 0.10\%. Our measurements correspond to more than an order of magnitude of improvement on previous realizations in Nd\(^{3+}\):YVO\(_4\) [18]. Since the material properties give the ultimate width of the spectral selection, they strongly influence the efficiency. This comparison shows that further improvements will demand material developments. This is especially true for the future backward configuration since Tm\(^{3+}\):YAG and Nd\(^{3+}\):YVO\(_4\) have only two ground states. This configuration does not allow, at this time, optical pumping in an auxiliary level and the application of control pulses [17].

We have developed an original spectral analysis. Our model directly predicts the observed efficiency for a given initial absorption spectrum that we measure independently. For a definite bandwidth and a fixed delay, the efficiency is limited by both the available optical thickness and the preparation method. In that case, due to the power broadening effect, the width and the height of the comb peaks are correlated. As a consequence, the efficiency and the spectral feature
of the protocol (delay and multi-mode capacity) are not independent. These characteristics are
critical for quantum repeater applications. A global estimation of the ultimate performances will
demand the development of original preparation methods.

Finally, we have shown that even for a very weak pulse, $\sim 0.5$ photon per pulse, the storage
itself should not add any noise to the incoming field. The efficiency and the noise feature then
render the AFC particularly attractive as a quantum memory protocol.

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