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Quantitative analysis of electrically detected Ramsey fringes in P-doped Si

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This work describes detection of the laser preparation and subsequent coherent manipulation of the quantum states of orbital levels of donors in doped Si, by measuring the voltage drop across an irradiated Si sample. This electrical signal, which arises from thermal ionization of excited orbital states, and which is detected on a millisecond time scale by a voltmeter, leads to much more sensitive detection than can be had using optical methods, but has not before been quantitatively described from first principles. We present here a unified theory which relates the voltage drop across the sample to the wave function of the excited donors, and compare its predictions to experiments in which pairs of picosecond pulses from the Dutch free-electron laser FELIX are used to resonantly and coherently excite P donors in Si. Although the voltage drop varies on a millisecond time scale we are able to measure Ramsey oscillation of the excitation on a picosecond time scale, thus confirming that the donor wave function, and not just its excited state population, is crucial in determining the electrical signal. We are also able to extract the recombination rate coefficient to the ground state of the donor as well as the photionization cross section from the excited state and phonon induced thermal ionization rate from the excited state. These quantities, which were previously of limited interest, are here shown to be important in the description of electrical detection, which, in our unoptimized configuration, is sensitive enough to enable us to detect the excitation of $\sim 10^7$ donors.

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I. INTRODUCTION

Previous experiments, using a free electron laser (FEL) source of coherent THz excitation, have shown that the orbital states of donors in Si can be manipulated in the same way as the orbital states of atoms or ions in traps [1,2]. However, phonon emission channels in Si lead to the rapid decay of excited states without a useful photon signal and this limits the number of excited donors which can be detected optically [3]. Very recently we have demonstrated the coherent creation and destruction, via Ramsey interference of picosecond pulse pairs, of orbital wave packets in Si:P with both optical and electrical read-out [4]. The electrical detection mechanism, termed photothermal ionization spectroscopy, is based on the much higher thermal ionization probability for an excited $2p_{\pm}$ state than the ground $1s$ state, implying that at finite temperature the sample conductivity directly reflects the strength of the $1s$ to $2p_{\pm}$ orbital transitions [4–7]. Electrical detection of orbital excitation is by orders of magnitude more sensitive than the optical technique (see also Refs. [8,9] for methods of electrical detection of spin resonances), enabling the detection of less than $10^5$ electrons excited into the conduction band. It has the advantage that the sample is also the detector, enabling a unified theory of excitation and detection, which links the slow electrical response (millisecond) directly to the picosecond dynamics of the laser excitation. In the present work we present such a quantitative analysis of earlier experiments, where we measure the voltage response of a current-biased P-doped Si sample under irradiation by pairs of FEL laser pulses, tuned to the THz frequency of the $1s\rightarrow 2p_{\pm}$ transitions. These pairs of pulses interact coherently to produce Ramsey interference fringes which we detect through variations in the sample conductance.

This two-pulse coherent excitation of a selected donor state is exactly the kind of manipulation envisaged in some quantum information processing schemes [10], so it is important for the modeling of such schemes to be sure that our understanding of the physics is complete. As we shall see below, the electron recombination rate from the conduction band, the photoionization cross section from the upper state and the thermal ionization rate from the upper state are parameters important in the full description of the excitation process. Fortunately...
we can extract these quantities from a comparison between theory and experiment, thus giving us a fully predictive model of the coherent manipulation of donors in Si by resonant laser excitation. To our knowledge, this is the first time that information has been obtained from both the absolute amplitude and the time dependence of coherent photoconductivity measurements.

In Sec. II we discuss the experiment and excitation dynamics of the donor electrons; in Sec. III we show, in detail, how the voltage drop across the sample is related to the population of the Si conduction band produced by the laser excitation, and how, by fitting the measured voltage drop as a function of time, we can extract not only the object of interest, the conduction band electron density, but also the free electron–ionized donor recombination coefficient and the effective circuit capacitance. In Sec. IV we compare the measured conduction band density as a function of laser interpulse separation with the result of theoretical predictions using a master equation, and find good agreement, thus confirming both the electrical and excitation models. These results, and what we learn from them, are discussed in Sec. V, and some conclusions are drawn in Sec. VI.

II. EXPERIMENT

All the experiments [11] are done as four-point, current-biased measurements with a phosphorus-doped silicon sample [12], doped at $2 \times 10^{20}$ donors m$^{-3}$ (see Fig. 1). Typically a constant current of 10 to 100 $\mu$A is passed between the outer terminals of the sample, and the voltage across the inner terminals under laser irradiation is recorded as a function of time on an oscilloscope. Our light source is the Dutch free-electron laser FELIX [13]. This produces trains of $\sim$100 transform-limited pulses with a duration of $\sim$5 ps, tuneable carrier frequency $\sim$9.5 THz, and interpulse interval of 40 ns. Each such train is called a macropulse, and each component of the train a micropulse. We use a beam splitter and delay stage to separate each micropulse in the train into a pulse pair with a separation of $t_d \sim$ ps. The micropulse energies of the two pulses in this pair, before final attenuation by the attenuator $A_2$ (see Fig. 1), are $E^{(1)} = 75 \mu$J and $E^{(2)} = 44 \mu$J, which are sufficiently similar for our purpose; the beams are focused to a waist of $0 = 9$ mm at the target. The final attenuations $A$ before being focused onto the sample with a parabolic mirror (PM).

The sample was mounted in vacuum on the cold finger of the exchanger at the top of the finger, was 9.5 K. We may calculate the thermal diffusion coefficient [15] for Si at 9.5 K is 0.69 m$^2$ s$^{-1}$ so the thermal relaxation time is $\sim \mu$s.

Thus there is no thermal recovery between micropulses, but complete thermal recovery between macropulses. The change in temperature during a macropulse has a negligible effect on conductance.

In the absence of FELIX we expect, and find, that the voltage is constant. If we irradiate the sample with a sequence of FELIX pulses the voltage changes with time. It drops when the FELIX pulses arrive, and then recovers to its original value. The details of the voltage drop depend upon the separation of the pulse pair and the laser intensity. This is because the FELIX pulses increase the conduction band electron density, thereby increasing the sample conductance, and thus reducing the voltage, which then recovers as recombination occurs.

The time scale for the measured voltage behavior is set by the electrical properties of the circuit—essentially $RC$,

\begin{equation}
\Delta T = \Delta I / C_s,
\end{equation}

where $\Delta I$ is the total energy absorbed by the sample over the course of a FELIX macropulse, and $C_s$ is the sample specific heat [14]. If we assume that all the FELIX energy goes into heating the sample we find

$\Delta T \sim 0.2$ K

for experiment 1, and correspondingly less for experiments 2 to 6. The thermal diffusion coefficient [15] for Si at 9.5 K is 0.69 m$^2$ s$^{-1}$ so the thermal relaxation time is $\sim \mu$s.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|}
\hline
$k$ & 1 & 2 & 3 & 4 & 5 & 6 \\
\hline
$A_2$ & 24.6 & 33.1 & 36.0 & 37.7 & 40.6 & 40.3 \\
\hline
\end{tabular}
\caption{The final attenuations: $A_2$ is the actual attenuation, in dB, used in the experiment labeled by index $k$.}
\end{table}
between macropulses, so we may assume complete relaxation of the sample between macropulses.

For each pulse pair separation, i.e., for each value of $t_d$, we record the voltage across the sample as a function of time. To avoid confusion with other time dependencies, we call the temporal parameter for these traces $t_{elec}$. Typically we record the traces for $t_{elec} \sim 300 \mu s$, very much longer than the times associated with the FELIX pulses.

Now we may calculate the sample resistance $R_s$ at conduction band density $n_e$ as [12]

$$R_s = \frac{L_e}{A_e n_e \mu},$$

where $L_e$ is the sample length, $A_e$ its cross sectional area, $\mu$ is the electron mobility, and $e$ is the charge on the electron. If the sample carries current $I_s$, the total voltage across the sample is

$$V_s = V_B + I_s R_s,$$

where $V_B$ is the barrier voltage, a function of $I_s$. Thus, knowing the $I$-$V$ characteristic of the sample, and hence $V_B$, the sample dimensions, and the electron mobility at the sample temperature, we may derive the equilibrium electron density $n_{eq}$ from $V_s$ and $I_s$.

Now we assume that this equilibrium density is maintained by some excitation process, with rate $S$, such as photoexcitation by black body radiation from the room temperature windows just above the sample, then the conduction band electron density $n_e$ satisfies

$$\dot{n}_e = -P n_e (n_c + n_e) + S (n_P - n_c - n_e),$$

where $n_P$ is the doping density, $n_c$ is the compensation density, and $P$ is the recombination rate coefficient. Noting that $n_e = 0$ when $n_e = n_{eq}$, we can eliminate $S$ and show that the conduction band density $n_e(t)$ as a function of time $t$ is

$$n_e(t) = \frac{n_0 \mu e^{-t/\tau_c} + n_{eq} (n_T + n_0)}{n_0 (1 - e^{-t/\tau_c}) + n_T},$$

where $n_0$ is the extra electron density added to the conduction band at time $t = 0$ by the FELIX pulses [16]. Here

$$n_e = \frac{(n_c + n_{eq})(n_P - n_c)}{n_P - n_c - n_{eq}}$$

is, approximately, the number of recombination centres at equilibrium

$$n_T = n_{eq} + n_e \sim 2n_{eq} + n_c$$

and finally

$$\tau_c = 1/(nP P)$$

is the characteristic recombination time [17]. Thus we can express the relaxation of the conduction band electron density to $n_{eq}$ in terms of the recombination rate coefficient $P$, the equilibrium conduction band electron density $n_{eq}$, the compensation density $n_c$, and the dopant density $n_P$.

We now assume that a micropulse pair arriving at $t_{elec} = 0$ promotes $n_1$ electrons to the conduction band. We can use Eq. (4) to calculate the conduction band electron density at $t_{elec} + 40 \mu s$, when the next micropulse pair arrives and adds its contribution, to give a total extra density in the conduction
hand. We can continue this process until all $M$ micropulse pairs have passed and the conduction band density is $n_M$ (see Fig. 3). On the time scale of the voltage change even the macropulse looks instantaneous, so the effect of a macropulse (see Fig. 3). On the time scale of the voltage change even the macropulse looks instantaneous, so the effect of a macropulse is to “instantaneously” promote $n_M$ electrons to the conduction band ($n_0 = n_M$). Equation (4) can then be used to calculate the subsequent conduction band density and, therefore, through Eq. (1), the sample resistance as a function of time.

III. ELECTRICAL MATTERS

We know $V_0$, the equilibrium voltage at current $I_s$, and from Eqs. (1), (2), and (4) we have an analytic form for the sample resistance as a function of time. Remarkably, for the circuit shown in Fig. 1, we may solve the Kirchhoff equations to obtain the voltage across the sample. It is

$$V(t) = V_B + (V_0 - V_B) \times \left( \Theta(t - t_p) \left\{ \frac{n_T + n_M}{n_T + n_M(1 - e^{-(t-t_p)/\tau_c})} \right\} \right)$$

$$\times \left\{ 2F_1 \left( \frac{\tau_c}{\tau_0} - \frac{n_T \tau_c}{n_0 \tau_0} ; 1 - \frac{\tau_c}{\tau_0} ; e^{-t/T_c} n_M \right) - e^{-t/T_c} 2F_1 \left( \frac{\tau_c}{\tau_0} - \frac{n_T \tau_c}{n_0 \tau_0} ; 1 - \frac{\tau_c}{\tau_0} ; n_M \frac{n_T + n_M}{n_T + n_M(1 - e^{-(t-t_p)/\tau_c})} \right) \right\}$$

$$+ e^{-t/T_c} \left( \frac{n_T}{n_T + n_M(1 - e^{-(t-t_p)/\tau_c})} \right)$$

$$+ \Theta(t_p - t) \right\},$$

(5)

where $t_p$ is the time at which the conduction band density increases by $n_M$, and

$$\tau_0 = C L_s/(\mu A_s n_0) \quad (=C(V_s - V_B)/I_s),$$

where $C$ is the circuit capacitance and $L_s/(\mu A_s n_0)$ is the sample equilibrium resistance, so that $\tau_0$ is the electrical time constant of the circuit. Here $L_s$ is the effective sample length, $A_s$ is the effective sample cross sectional, and $\mu$ is the electron mobility. We use $L_s = 2 \text{ mm}$, $A_s = 1.6 \times 10^{-6} \text{ m}^2$, derived from the sample dimensions, and $\mu = 10 \text{ m}^2 \text{ V}^{-1} \text{ s}^{-1}$, derived from a fit to literature values [18]. In Eq. (5) $2F_1(a,b;c;z)$ is a hypergeometric function and $\Theta(t)$ is a Heaviside step function. We also use $t$ as a general time variable; it is equivalent to $t_{\text{elec}}$ in the experiments.

There are three unknowns in Eq. (5): the circuit capacitance $C$, the recombination coefficient $P$, and the conduction band density $n_M$. These can be obtained by fitting Eq. (5) to the experimental data—the voltage traces $V(t_{\text{elec}})$ for each pulse separation $t_d$. We expect the recombination coefficient $P$, and the circuit capacitance $C$ to be independent of the laser pulse area [19] or pulse delay $t_d$, but the density of electrons excited to the continuum $n_M$ to be a strong function of both these parameters.

We focus first on experiments 2 to 6. These were done at high attenuation and, as we shall see, in contrast to experiment 1, the voltage drops are reasonably small; saturation is not a problem and fitting Eq. (5) to the experimental voltage profiles is straightforward. In Fig. 4 we show typical fits of the voltage drop to Eq. (5), for experiments 2 to 6, whence we find [20]

$$P = (5.3 \pm 0.7) \times 10^{-13} \text{ m}^3 \text{ s}^{-1}$$

and

$$C = 1.24 \pm 0.14 \text{ nF}.$$
FIG. 5. (Color online) The conduction band density as a function of interpulse delay $t_d$ for experiment 5. The actual experimental values are shown as a thin (blue) line. The upper border (shown in red) and the lower border (shown in green) are used in the fitting procedure described below. Also shown, as an inset, is the experimental power spectrum, normalized to unit height, as a function of frequency in THz (blue points), and in black, a Gaussian fit to it, showing that the rapid oscillations are indeed at the $1_s$-$2p_x$ transition frequency of 9.49 THz.

strongly oscillating function of $t_d$, with a smooth envelope [21]. Our interest, however, is not the density in itself, but the coherent excitation of donor bound states by intense laser pulses [10,22]. As we have described in the Introduction, and illustrated in Fig. 2, we believe that the dominant mechanism for population of the conduction band is thermal ionization from the upper level excited by the FELIX pulses: this implies that the conduction band density is proportional to the upper state population produced by the laser excitation. Thus, the conduction band density serves as a measure of the excitation of the bound state by a FELIX pulse pair. All the processes shown in Fig. 2 can be included in a master equation [23]. If the solutions of this master equation reproduce the experimental results, we shall have confirmation that the electrical detection technique described here is capable of providing a useful alternative to the optical methods more usually used to detect coherent bound-state excitation.

From previous measurements using standard optical and electrically detected Fourier-transform infrared (FTIR) spectroscopy at different sample temperatures [24], we have good values for $|d|$, the radial matrix element between the ground and excited donor states, and for $G$, the thermal ionization rate from the upper level. $T_1$ and $T_2$, the population and coherence decay rates, are known from pump-probe [2] and echo [1] experiments, and, of course $\tau_{\text{pulse}}$, the laser pulse duration is obtained from standard pulse autocorrelation measurements. Thus, these parameters are known independent of the current experiments, and their values are shown in Table II.

IV. MODELING THE EXCITATION

Figure 2 shows all the important processes included in our model of the excitation process, which, of course, can be described by a master equation (in the rotating wave approximation [25])

$$
\begin{pmatrix}
\dot{\rho}_{11} \\
\dot{\rho}_{12} \\
\dot{\rho}_{21} \\
\dot{\rho}_{22}
\end{pmatrix}
= \begin{pmatrix}
0 & -i\frac{\Omega_2(t)}{2} & 0 & -i\frac{\Omega_2(t)}{2} \\
-i\frac{\Omega_2(t)}{2} & -i\frac{\Omega_1(t)}{2} & -G & i\frac{\Omega_2(t)}{2} \\
0 & -G & -i\frac{\Omega_2(t)}{2} & 0 \\
-i\frac{\Omega_2(t)}{2} & i\frac{\Omega_2(t)}{2} & -G & -i\frac{\Omega_1(t)}{2}
\end{pmatrix}
\begin{pmatrix}
\rho_{11} \\
\rho_{12} \\
\rho_{21} \\
\rho_{22}
\end{pmatrix}
$$

(6)

where $\rho$ is the density matrix describing the state of a donor, and, as indicated in Fig. 2, the other parameters have the following meaning: $\Omega_2(t) = \Omega_2^*(t) = e d \cdot E(t)/\hbar$ is the instantaneous Rabi frequency produced by the laser pulses. Here $E(t)$ is the instantaneous laser field envelope and $e d$ is the transition dipole matrix element; $\Delta$ is the detuning between the laser carrier frequency and the donor transition frequency, $\gamma_21$ is the population decay rate from level 2 to level 1, and $\Gamma$ is the dephasing rate between levels 1 and 2 (i.e., $\gamma_21 = 1/T_1$ and $\Gamma = 1/T_2$), $\sigma_2$ is the photoionization cross section from level 2, $F(t) = I(t)/\epsilon$ is the photon flux produced by the laser, where $I(t)$ is the instantaneous laser intensity and $\epsilon$ is the laser photon energy, and $G$ is the thermal photoionization rate from the upper level. Of course $\Omega_2(t)$ and $F(t)$ are functions both of time and $r$, the radial distance from the beam axis, reflecting the two-pulse nature of the excitation and the Gaussian beam profile, respectively. Finally, we have

$$
n_1(t,t_d) = n_p \int_{-\infty}^{\infty} d\Delta g(\Delta) \int_0^R 2\pi r dr n_{\text{ion}}(t,t_d,\Delta,r)/S,
$$

(7)

with

$$
n_{\text{ion}}(t,t_d,\Delta,r) = [1 - \rho_{11}(t,t_d,\Delta,r) - \rho_{22}(t,t_d,\Delta,r)],
$$

where $g(\Delta)$ is the inhomogeneous broadening profile, $R (= 1$ mm) is the radius of the illuminated region, set by an aperture in front of the sample, and $S \sim \pi R^2$ is the effective area occupied by the conduction band electrons.

Thus, as Eq. (6) shows, in addition to the parameters in Table II we need to know the inhomogeneous broadening of the sample $g(\Delta)$ and the photoionization cross section from the upper level $\sigma_2$. We consider two inhomogeneous profiles,

---

### Table II. The known data

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$</td>
<td>d</td>
</tr>
<tr>
<td>$T_1$ (ps)</td>
<td>200</td>
</tr>
<tr>
<td>$T_2$ (ps)</td>
<td>160</td>
</tr>
<tr>
<td>$G$ ($s^{-1}$)</td>
<td>$2.5 \times 10^6$</td>
</tr>
<tr>
<td>$\tau_{\text{pulse}}$ (ps)</td>
<td>3.7</td>
</tr>
</tbody>
</table>

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165310-5
 Fit | $f$ (%) | $T_{\text{eff}}$ | $W$ (THz) | $\sigma_2$ ($10^{-20}$ m$^2$) |
<table>
<thead>
<tr>
<th></th>
<th></th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>L</td>
<td>2.3 ± 0.8</td>
<td>0.95 ± 0.1</td>
<td>0.057 ± 0.01</td>
<td>2.5 ± 1</td>
</tr>
<tr>
<td>G</td>
<td>2.5 ± 0.8</td>
<td>0.83 ± 0.08</td>
<td>0.073 ± 0.01</td>
<td>2.6 ± 0.4</td>
</tr>
</tbody>
</table>

In summary, the fitting parameters $f$ is the percentage of third harmonic energy in the FELIX pulse, $T_{\text{eff}}$ is the effective electric field transmission coefficient into the cryostat, $W$ is the full width half maximum (FWHM) value of the inhomogeneous broadening, and $\sigma_2$ is the photoionization cross section from the upper level. L assumes a Lorentzian and G a Gaussian shape for the inhomogeneous broadening. The Lorentzian fit is marginally better.

### Table III: The fitting parameters

<table>
<thead>
<tr>
<th>Fit</th>
<th>$f$ (%)</th>
<th>$T_{\text{eff}}$</th>
<th>$W$ (THz)</th>
<th>$\sigma_2$ ($10^{-20}$ m$^2$)</th>
</tr>
</thead>
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<td>2.6 ± 0.4</td>
</tr>
</tbody>
</table>

Finally, there are two purely experimental features which we must take into account. First, absolute intensity measurements in this wavelength region are very uncertain, and in our setup it is difficult to know how much transmission loss we have between the target and the point in front of the polypropylene cryostat window where the beam intensity was measured, with a commercial portable pyroelectric power meter. We account for this by introducing $T_{\text{eff}}$, defined so that

$$E_{\text{target}} = T_{\text{eff}} E_{\text{meas}},$$

where $E_{\text{meas}}$ is the laser envelope field derived from the measured laser power, and $E_{\text{target}}$ is the laser field at the surface of the Si target. Previous experience [1] suggests that $T_{\text{eff}} \sim 0.7$. Second, the FELIX radiation has a small third harmonic component [26]. Although this is at the few percent level it is important in these experiments because third harmonic photons are sufficiently energetic to ionize ground state donors directly. Since direct ionization does not show Ramsey fringes, its effect is to add an independent background to the results calculated from Eq. (6). This background is of magnitude

$$n_s \sim f n_p A \sigma_{\text{1s}}(3 \epsilon_{2\text{pl}}) \frac{\mathcal{E}_{\text{in}}}{3\pi\epsilon_0^2 \epsilon_{2\text{pl}}},$$

where $f$ is the fraction of the third harmonic, $\epsilon_{2\text{pl}}$ is the transition energy, $\sigma_{\text{1s}}(3 \epsilon_{2\text{pl}})$ is the $1s_\text{A}$ ionization cross section at the third harmonic energy ($\sim 3 \times 10^{-20}$ m$^2$ from hydrogenic scaling), $A$ is the beam attenuation, and $\mathcal{E}_{\text{in}}$ is the total beam energy at the target.

The important physical content of the conduction band density is the envelope of the rapidly oscillating signal. We make a fit of the theoretical expression derived from Eqs. (7) and (8) to the results from experiments 2 to 6 simultaneously using the beam transmission $T_{\text{eff}}$, inhomogeneous width $W_L$ or $W_G$, cross section $\sigma_2$, and third harmonic component $f$ as fitting parameters. A typical result is shown in Fig. 6, and the parameters given in Table III. The errors reflect the spread of values over the individual experiments. The fits themselves show errors of $\sim 13\%$ per point, averaged over experiments 2 to 6. Nominally the error in the Lorentzian fit is smaller—13.1% compared with 13.3%—but this difference is, of course, not significant. Finally, we may remark that we have checked that both the fits to the voltage drops and the fits to the fringe envelopes are robust against variations in their starting parameters.

### V. DISCUSSION

Electrically detected FTIR spectroscopy has been performed on our sample, and both Lorentzian and Gaussian fits made to the $2P_{\perp}$ line. It is found that the Lorentzian shape gives a systematically better fit [27]. The results are that, for a Lorentzian fit, the full width half maximum is

$$W_{\text{L}}^{\text{FTIR}} = 0.041 \pm 0.007 \text{ THz},$$

and for a Gaussian fit, the FWHM is

$$W_{\text{G}}^{\text{FTIR}} = 0.047 \pm 0.01 \text{ THz}.$$

Thus the values for the widths given in Table III are in broad agreement with the FTIR measurements, particularly as the
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FIG. 7. (Color online) The analog of Fig. 4 for experiment 1. Here the voltage drop is typically much larger than in the previous cases.

linewidth is sensitive to strain induced by the sample mounting method. However, the agreement between the widths derived from the Ramsey fringe experiments and the FTIR spectra is better for the Lorentzian line shape than for the Gaussian. Once again, the difference is not conclusive (a 1.3 standard deviation difference for the Lorentzian shape as opposed to a 1.8 standard deviation difference for the Gaussian) but, as before, the Lorentzian shape is preferred. The photoionization cross section from the upper state has a large error but is consistent with the value expected from hydrogenic scaling. Furthermore, \( T_{\text{eff}} \) appears to be well determined by the fits, and suggests that 70%–90% of the measured input power is transferred to the surface of the sample. The value of \( f \), the fraction of total energy in the third harmonic, is also at the level expected.

We now consider experiment 1. In Fig. 7 we show the analog of Fig. 4 for experiment 1, the fits of the voltage response as a function of \( t_{\text{elec}} \). The circuit capacitance obtained from these fits, the value which minimizes the maximum error over the set of voltage drops is

\[
C_1 = 2.6 \, \text{nF}.
\]

Notice first that the fits are somewhat poorer, but, more important, because of the more intense excitation, and consequent increase in conduction band occupancy, the sample resistance is lower and thus, from Eq. (2), the total voltage drop is larger than in the other experiments. Now Eq. (2) also shows that the voltage drop must saturate—it can be no larger than \( V_0 - V_B \). This has an important consequence that we can deduce from Eq. (5). For a given set of parameters \( n_M, \tau_0, \tau_C, \) etc., we may calculate \( \Delta V = V_0 - V_{\text{min}} \), where \( V_{\text{min}} \) is the minimum voltage attained for the chosen set of parameters, i.e., \( \Delta V \) is the maximum voltage drop attained for these parameters. We assume \( \tau_0, \tau_C, \) etc., are fixed, and focus on the behavior as a function of \( n_M \). We then have

\[
\Delta V = F(n_M),
\]

where \( F \) is a function which could be derived from Eq. (5). We may invert this equation and, taking \( G = F^{-1} \), derive

\[
\frac{\delta n_M}{n_M} = h(\Delta V)\frac{\delta \Delta V}{\Delta V},
\]

with

\[
h(\Delta V) = \frac{\Delta V G(\Delta V)/G(\Delta V)}{\Delta V}
\]

(10)

(of course \( h \) depends upon \( \tau_0, \tau_C, \) etc., as well). We may interpret Eq. (9) as meaning that the relative uncertainty in the continuum electron density \( \delta n_M/n_M \) is proportional to the relative uncertainty in the maximum voltage drop \( \Delta V/\Delta V \) with the constant of proportionality given by \( h(\Delta V) \). Now, for each experiment we have ~1500 values of \( \Delta V \), one for each \( t_d \) from which we may construct a probability density function

\[
P_k(\Delta V) d\Delta V,
\]

which shows the distribution [28] of maximum voltage drops for experiment \( k \). Thus, \( P_k \) defines the range over which \( \Delta V \) varies for experiment \( k \) and \( h(\Delta V) \) shows how errors in \( \Delta V \) are propagated into \( n_M \). We show both \( h(\Delta V) \) and \( P_k(\Delta V) \) in Fig. 8. The implication of this figure is that even if the experimental errors in \( V(t_{\text{elec}}) \) are the same for all the experiments, nevertheless the uncertainty in extracting the continuum electron density \( n_M \) from the behavior of \( V(t_{\text{elec}}) \) grows with increasing maximum voltage drop, i.e., with increasing laser intensity.

In Fig. 9 we compare the experimental continuum density envelope with the theoretical prediction. Since these experimental results were not used in obtaining the quantities needed by theory (see Table III), this continuum density envelope is a true prediction. It agrees with experiment less well than in the other cases: we attribute this to the difficulty of extracting the continuum electron densities adumbrated above. Nevertheless, we believe the agreement is acceptable.
VI. CONCLUSIONS

In summary, we have developed a theory of the coherent excitation of donors in Si and applied it to a description of the electrical detection of the coherent response of P donors to irradiation by pairs of THz laser pulses. Electrical detection is sensitive. Typically, in these experiments we detect the ionization of about one donor in $10^4$, within a volume of $\sim 3 \times 10^{-9} \text{ m}^3$ implying a total of $\sim 10^7$ excited electrons, and, as Fig. 6 shows, the limit is certainly an order of magnitude better than this. However, as the analysis presented above shows, the small number of conduction band electrons is a result of the low thermal ionization rate of the excited state, a condition exacerbated by the use of lower operating temperatures. To exploit the method fully, direct readout methods which selectively ionize the excited state at low temperatures (where the photothermal channel is shut down), for example by using a long-wavelength intense laser, will have to be developed. These methods will then be suitable for dilute samples, where the saturation shown in Fig. 8 is not a problem.

However, what is more important is that in this study we have identified all the important physical processes which must be included in a theory of the coherent excitation of donors in Si. Many of these are the result of the solid-state environment of the donors, and would be absent in a trapped ion or atom. Nevertheless, their inclusion in the master equation that describes the excitation is straightforward, so that essentially the same techniques used in the description of the coherent control of ions and atoms in traps can be applied to donors in Si, once again suggesting that Si provides a “vacuumlike” environment for donors [1,2].

However, the analogy is not perfect. A major difference is that electrical detection, as used here, is sensitive to recombination, thus focusing attention on the effective recombination rate coefficient $P$. Here we find

$$P = 5.3 \times 10^{-13} \text{ m}^3 \text{ s}^{-1},$$

which is close to

$$P_{BR} = 2 \times 10^{-13} \text{ m}^3 \text{ s}^{-1},$$

the value estimated from the Brown and Rodriguez [20] value of recombination cross section. However, Norton et al. [29] give the much larger value of

$$P_N = 2.5 \times 10^{-11} \text{ m}^3 \text{ s}^{-1}$$

close to the value

$$P_B = 1.3 \times 10^{-11} \text{ m}^3 \text{ s}^{-1}$$

found by Bowyer et al. [7]. It should be noted that the sample used here is similar to those considered by Brown and Rodriguez and Norton et al., but very different from the metal oxide semiconductor field effect transistor (MOSFET) used by Bowyer et al. Recombination is a complex cascade process and the effective recombination rate coefficient will depend at least on temperature, donor density, trap density (which will vary strongly between bulk silicon and MOSFETs), and sample phonon spectrum, so the difference between our value and that in Ref. [7] is not surprising.

In summary, the important conclusion is that techniques which have been successfully applied to the modeling of quantum information processing in trapped ions and atoms [30] can, with little modification, also be used to model similar processes in donors in Si.

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