

Electrical Spectroscopy of Intervalley Relaxation in WSe₂ Transistors

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We show that the transconductance of multilayer WSe₂ field-effect transistors serves as a direct electrical spectrometer of the intervalley relaxation time τ_{iv} , previously accessed mainly by ultrafast optical techniques. Extending an equilibrium valley-thermodynamics framework with a single relaxation equation for the Γ -valley carrier fraction $f_{\Gamma}(t)$, we predict three signatures: (i) a Lorentzian transconductance $g_m(\omega) = g_{m,0} + g_{m,v}^0/(1 + i\omega\tau_{iv})$, whose imaginary part peaks at $\omega_c = \tau_{iv}^{-1}$ with opposite signs for bilayer and trilayer; (ii) a two-stage current transient after a gate step, exhibiting bilayer overshoot or trilayer undershoot; and (iii) sweep-rate-proportional hysteresis whose gate-voltage profile and the layer-dependent transconductance sign reversal distinguish valley from trap-induced dynamics. All three signatures provide quantitative electrical access to τ_{iv} with standard radio-frequency (rf) and direct-current (dc) instrumentation.

I. INTRODUCTION

A transistor channel hosting multiple valleys with different effective masses carries a slow internal degree of freedom: the intervalley carrier distribution. When the intervalley equilibration time τ_{iv} is not negligibly short compared with the gate-modulation period, this distribution cannot follow the gate adiabatically, imprinting itself on all time-dependent transport observables. The consequences form a recognizable set of dynamic signatures — frequency-dependent transconductance, two-stage current transients, and sweep-rate-proportional hysteresis — whose gate-voltage profiles and the layer-dependent transconductance sign reversal distinguish valley dynamics from charge-trapping effects. Accessing τ_{iv} through standard rf and dc instrumentation would provide a direct spectroscopic route to an internal scattering channel previously accessed mainly through ultrafast optical spectroscopy.

Multilayer WSe₂ provides a layer-tunable platform for the study of these effects [1–8]. The valence-band maximum shifts from the K point in the monolayer toward the Γ point with increasing layer number, driven by interlayer coupling and spin-orbit interaction [9–13], placing $\Delta_{K\Gamma} \approx k_B T$ near room temperature in the bilayer. This competition enables gate-tunable redistribution between the K and Γ valleys, producing a valley-dependent transconductance contribution $g_{m,v} = -\Delta\mu(W/L)V_{DS}qp\chi_v$ parameterized by the valley susceptibility $\chi_v \equiv \partial f_{\Gamma}/\partial V_{GS}$, completing the decomposition $g_m = g_{m,0} + g_{m,v}$ [14–19]. Such treatments assumed quasi-static valley equilibrium. The relaxation time τ_{iv} in WSe₂ is set by phonon-mediated scattering at room temperature (sub-ps to a few ps) [20–26] but can be substantially longer at reduced temperatures or under gate-induced non-equilibrium conditions [27–29], potentially

placing $\omega_c = \tau_{iv}^{-1}$ in the accessible MHz–GHz window. To date, τ_{iv} in WSe₂ has been probed almost exclusively by ultrafast optical spectroscopy [30, 31], which requires specialized laser and optical cryostat infrastructure beyond conventional electrical transport setups.

In this paper, we extend the equilibrium valley framework to the non-equilibrium regime through a single relaxation equation for $f_{\Gamma}(t)$, characterized by the intervalley relaxation time τ_{iv} . This minimal extension predicts three measurable consequences of delayed intervalley redistribution: (i) a Lorentzian frequency dependence of $g_m(\omega)$ with characteristic frequency $\omega_c = \tau_{iv}^{-1}$; (ii) a two-stage current transient after a gate step exhibiting bilayer overshoot or trilayer undershoot, determined by the sign of $g_{m,v}^0$; and (iii) sweep-rate-proportional hysteresis whose gate-voltage profile follows $qp\chi_v(V_{GS})$ and reverses sign between the two layer numbers. In each case, the transconductance sign reversal with layer number and the gate-voltage dependence allow valley-origin dynamics to be distinguished from charge-trapping effects, and together they provide a quantitative electrical route to τ_{iv} . Figure 1(a) shows the device geometry and measurement configuration, Fig. 1(b) illustrates the resulting valley-population lag, and Fig. 1(c) displays the predicted spectroscopic signature in $\text{Im}[g_m(\omega)]$.

II. DYNAMIC VALLEY MODEL

The equilibrium valley-thermodynamics framework [15] yields a transconductance decomposition $g_m^{(0)} = g_{m,0} + g_{m,v}^0$, where $g_{m,0}$ is the conventional charge-accumulation term and

$$g_{m,v}^0 = -\frac{W}{L}V_{DS}\Delta\mu qp\chi_v, \quad (1)$$

with $\Delta\mu = \mu_K - \mu_{\Gamma} > 0$, q the elementary charge, p the total hole density, and $\chi_v \equiv \partial f_{\Gamma}/\partial V_{GS}$ the valley susceptibility. The sign of $g_{m,v}^0$ follows that of $\Delta_{K\Gamma}$: negative for the bilayer, positive for the trilayer.

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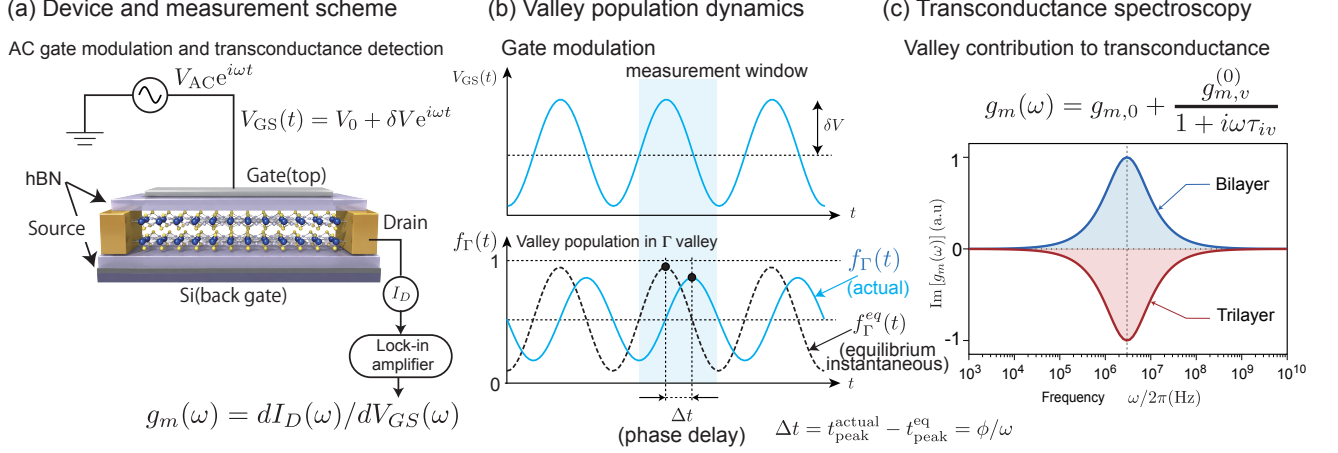


FIG. 1. Overview of transconductance spectroscopy of intervalley relaxation. (a) Schematic of a multilayer WSe₂ transistor in dual-gate geometry with hexagonal boron nitride (hBN) dielectric [32, 33]. An alternating-current (ac) gate voltage $V_{GS}(t) = V_0 + \delta V e^{i\omega t}$ drives the device; the drain current $I_D(\omega)$ is detected by a lock-in amplifier to extract $g_m(\omega) = dI_D(\omega)/dV_{GS}(\omega)$. (b) Valley population $f_\Gamma(t)$ in the Γ valley (solid) compared to the instantaneous equilibrium value (dashed), illustrating the phase delay $\Delta t = \phi/\omega$ arising from the finite intervalley relaxation time τ_{iv} . (c) Predicted imaginary part of the valley transconductance $\text{Im}[g_m(\omega)]$ for bilayer (blue) and trilayer (red) WSe₂. Opposite-sign Lorentzian peaks centered at $\omega_c = \tau_{iv}^{-1}$ provide a direct spectroscopic readout of τ_{iv} .

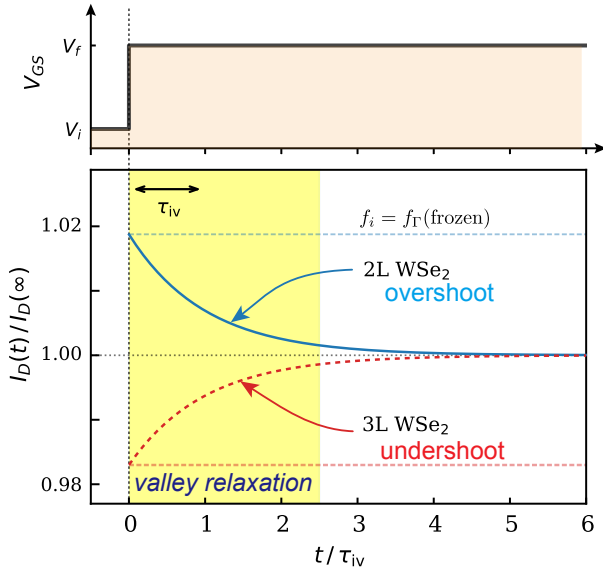


FIG. 2. Two-stage current response to a gate step $\Delta V_{GS} = 0.2$ V (upper panel) for bilayer (blue) and trilayer (red) WSe₂ (lower panel, thick lines). Thin dashed lines ($f_\Gamma = f_i$, frozen) show the valley-frozen reference: charge responds instantaneously ($\tau_{RC} \ll \tau_{iv}$) while f_Γ remains at its pre-step value f_i . Valley relaxation then drives the current toward $I_D(\infty)$: overshoot for the bilayer and undershoot for the trilayer, with a characteristic timescale τ_{iv} . The opposite-sign departures reflect $\text{sgn}(g_{m,v}^0)$ reversing between layers.

The equilibrium fraction of holes occupying the Γ val-

ley at gate voltage V_{GS} is

$$f_\Gamma^{\text{eq}}(V_{GS}) = \frac{p_\Gamma^{\text{eq}}(V_{GS})}{p_K^{\text{eq}}(V_{GS}) + p_\Gamma^{\text{eq}}(V_{GS})}, \quad (2)$$

where p_ν^{eq} is the equilibrium two-dimensional hole density in valley $\nu \in \{K, \Gamma\}$, determined from Fermi-Dirac statistics with valley density of states $D_\nu \propto m_\nu^*$ and the gate-controlled Fermi energy [15]. The susceptibility χ_ν peaks near the K - Γ band-crossing condition and vanishes in both the subthreshold ($p \rightarrow 0$) and heavily accumulated limits, confining the valley contribution $g_{m,v}^0 \neq 0$ to the above-threshold partially occupied regime.

Layer-dependent sign reversal.— The sign of χ_ν —and hence of $g_{m,v}^0$ —is determined by which valley constitutes the valence-band maximum, a property that reverses sign with layer number in WSe₂ [9, 11]. In the bilayer ($\Delta_{K\Gamma} = +26$ meV), the K valley lies at higher hole energy: holes at low carrier density reside predominantly in K , and a positive gate step drives them progressively into Γ , giving $\chi_\nu > 0$ and hence $g_{m,v}^0 < 0$. In the trilayer ($\Delta_{K\Gamma} = -49$ meV), the ordering reverses: holes first fill Γ , and the Γ fraction decreases as K also becomes occupied with increasing gate voltage, giving $\chi_\nu < 0$ and $g_{m,v}^0 > 0$. This layer-dependent transconductance sign reversal is absent in single-valley materials or in systems where $\Delta_{K\Gamma}$ has a fixed sign, and is specific to the layer-dependent band structure of multilayer WSe₂.

We extend this to the non-equilibrium regime with a single relaxation equation for the Γ -valley fraction $f_\Gamma(t) = p_\Gamma(t)/p(t)$:

$$\frac{df_\Gamma}{dt} = \frac{f_\Gamma^{\text{eq}}[V_{GS}(t)] - f_\Gamma(t)}{\tau_{iv}}, \quad (3)$$

the transistor analogue of Debye relaxation in dielectrics [34]. Extensions to distributed τ_{iv} (Cole-Cole, stretched exponential) follow directly. The single- τ_{iv} case is the falsifiable baseline.

Separation of timescales.— Equation (3) is valid in the regime $\tau_{RC} \ll \tau_{iv}$, where the charge relaxation time $\tau_{RC} = R_{ch}C_{ox}$ is much shorter than the intervalley relaxation time. Under this assumption, the total hole density $p(t) = p[V_{GS}(t)]$ tracks V_{GS} instantaneously, while $f_{\Gamma}(t)$ lags. The source and drain contacts are assumed to act as thermal reservoirs that define the carrier injection distribution at the channel boundaries, while transport within the channel is treated in the low-bias diffusive limit [35, 36]. The channel hole density $p(V_{GS}) = C_{ox}(V_{GS} - V_{th})/q$, with V_{th} the threshold voltage, responds electrostatically to the gate. The intervalley fraction $f_{\Gamma}(t)$ modulates the effective channel mobility

$$\mu_{\text{eff}} = \mu_K(1 - f_{\Gamma}) + \mu_{\Gamma}f_{\Gamma} = \mu_K - \Delta\mu f_{\Gamma}, \quad (4)$$

but leaves p unchanged by intervalley redistribution. This distinguishes the present mechanism from charge-trapping memtransistors, in which the trapped charge renormalizes the channel band edge and gate efficiency [35]. In the present model, no additional interface capacitance or band-edge shift is introduced. The dynamic drain current is then

$$I_D(t) = \frac{W}{L} \mu_{\text{eff}}[f_{\Gamma}(t)] qp[V_{GS}(t)] V_{DS}. \quad (5)$$

Figure 2 illustrates the physical picture: a gate-voltage step ($V_{GS} : V_1 \rightarrow V_2$) causes the charge to jump instantaneously while f_{Γ} exponentially relaxes toward the new equilibrium, producing a two-stage current response.

Table I lists the numerical parameters used throughout this work. The effective masses and K - Γ splittings are taken from first-principles calculations [11, 13], and the mobilities are representative experimental values [6, 7]. The intervalley relaxation time τ_{iv} enters only as a free parameter; no specific numerical value is assumed in any figure.

III. FREQUENCY-DEPENDENT TRANSCONDUCTANCE

For a small ac modulation $V_{GS}(t) = V_0 + \delta V e^{i\omega t}$, linearizing Eq. (3) gives the *dynamic valley susceptibility*

$$\chi_v(\omega; V_0) \equiv \frac{\delta f_{\Gamma}}{\delta V} = \frac{\chi_v(V_0)}{1 + i\omega\tau_{iv}}, \quad (6)$$

a Lorentzian whose pole is at $\omega_c = \tau_{iv}^{-1}$. Substituting into the linearized current of Eq. (5):

$$g_m(\omega) = g_{m,0}(V_0) + \frac{g_{m,v}^0(V_0)}{1 + i\omega\tau_{iv}}. \quad (7)$$

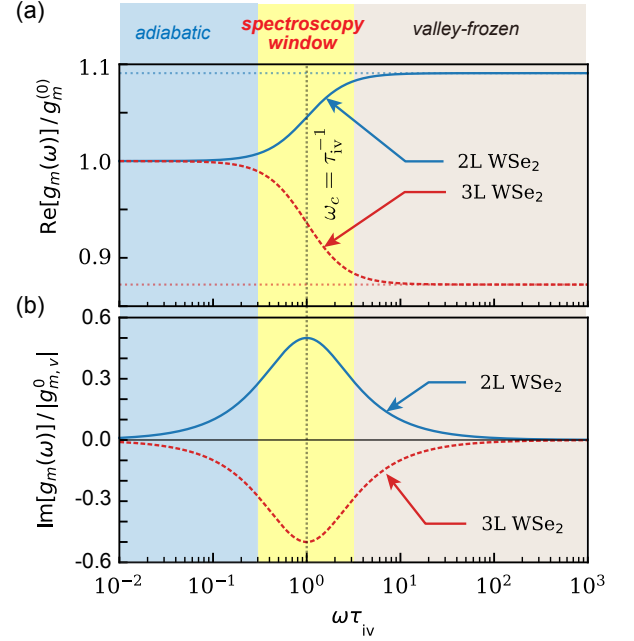


FIG. 3. Frequency-dependent transconductance for bilayer (solid, blue) and trilayer (dashed, red) WSe₂ at $T = 300$ K, evaluated at the gate voltage where $|g_{m,v}^0|$ is maximum. Shaded regions mark three dynamical regimes: adiabatic ($\omega\tau_{iv} \ll 1$), spectroscopy window ($\omega\tau_{iv} \sim 1$, yellow), and valley-frozen ($\omega\tau_{iv} \gg 1$). (a) Real part $\text{Re}[g_m(\omega)]/g_m^0$ normalized to the dc value; dotted lines are the high-frequency limits $g_{m,0}/g_m^0$. (b) Imaginary part $\text{Im}[g_m(\omega)]/|g_{m,v}^0|$; the extremum at $\omega_c = \tau_{iv}^{-1}$ (vertical dotted) has opposite sign for bilayer ($g_{m,v}^0 < 0$) and trilayer ($g_{m,v}^0 > 0$).

The charge-accumulation term $g_{m,0}$ is frequency-independent (instantaneous charge relaxation); all dispersion is carried by the valley term. Equation (7) is the central result of this paper.

Separating real and imaginary parts:

$$\text{Re}[g_m(\omega)] = g_{m,0} + \frac{g_{m,v}^0}{1 + (\omega\tau_{iv})^2}, \quad (8)$$

$$\text{Im}[g_m(\omega)] = -\frac{g_{m,v}^0 \omega\tau_{iv}}{1 + (\omega\tau_{iv})^2}. \quad (9)$$

Three limiting behaviors follow immediately: (i) at $\omega\tau_{iv} \ll 1$, $\text{Re}[g_m] = g_m^0$ recovers the equilibrium anomaly of Ref. 15; (ii) at $\omega\tau_{iv} \gg 1$, $\text{Re}[g_m] \rightarrow g_{m,0}$, the anomaly is completely suppressed; (iii) at $\omega = \omega_c$, $|\text{Im}[g_m]|$ reaches its maximum of $|g_{m,v}^0|/2$, providing a sharp spectroscopic feature that directly identifies τ_{iv} . Because $g_{m,v}^0 < 0$ for the bilayer and $g_{m,v}^0 > 0$ for the trilayer, the imaginary part peaks with *opposite signs* for the two systems, providing an additional layer-number signature.

These Kramers-Kronig-related dispersive and absorptive components are shown in Figs. 3(a) and 3(b), respectively. The high-frequency limit $g_{m,0}/g_m^0$ exceeds

TABLE I. Parameters used in the two-valley model.

Quantity	Value	Description
m_K^*	$0.40 m_0$	K -valley effective mass
m_Γ^*	$1.00 m_0$	Γ -valley effective mass
$\Delta_{K\Gamma}(2L)$	$+26 \text{ meV}$	K - Γ valley splitting, bilayer
$\Delta_{K\Gamma}(3L)$	-49 meV	K - Γ valley splitting, trilayer
μ_K	$100 \text{ cm}^2/(\text{V}\cdot\text{s})$	K -valley hole mobility
μ_Γ	$30 \text{ cm}^2/(\text{V}\cdot\text{s})$	Γ -valley hole mobility
C_{ox}	49 mF/m^2	Gate oxide capacitance per unit area ^a
$W = L$	$1 \mu\text{m}$	Channel width and length
V_{DS}	50 mV	Drain-source bias voltage
T	300 K	Temperature

^a Equivalent oxide thickness 0.7 nm.

unity for the bilayer (the valley suppression disappears) and falls below unity for the trilayer (the valley enhancement disappears), directly reflecting the sign of $g_{m,v}^0$ [Fig. 3(a)]. A frequency sweep at fixed $V_{GS} = V_{GS}^*$ (the peak- $|g_{m,v}^0|$ voltage) yields ω_c from either the -3 dB rolloff of $\text{Re}[g_m] - g_{m,0}$ or the peak of $|\text{Im}[g_m]|$, with both criteria giving the same τ_{iv} [Fig. 3(b)].

IV. TRANSIENT RESPONSE AND HYSTERESIS

Step response.— For a step $V_{GS} : V_i \rightarrow V_f$ at $t = 0$, the charge density jumps to $p(V_f)$ instantaneously while f_Γ satisfies Eq. (3) with initial condition $f_\Gamma(0^-) = f_\Gamma^{\text{eq}}(V_i) \equiv f_i$. The solution is $f_\Gamma(t) = f_f + (f_i - f_f)e^{-t/\tau_{iv}}$ with $f_f \equiv f_\Gamma^{\text{eq}}(V_f)$, giving

$$I_D(t) = I_D(\infty) + \Delta I_D^{\text{tr}} e^{-t/\tau_{iv}}, \quad (10)$$

where $I_D(\infty)$ is the equilibrium current at V_f and the transient amplitude is

$$\frac{\Delta I_D^{\text{tr}}}{I_D(\infty)} = \frac{\Delta\mu}{\mu_{\text{eff}}(V_f)} (f_f - f_i). \quad (11)$$

The sign of ΔI_D^{tr} is governed by the layer number: for bilayer WSe_2 , a positive gate step drives holes from the light K toward the heavy Γ valley ($f_f > f_i$), so $\Delta I_D^{\text{tr}} > 0$ and the current *overshoots* before decaying toward equilibrium. For trilayer, the reverse transfer gives $f_f < f_i$, $\Delta I_D^{\text{tr}} < 0$, and the current *undershoots* before recovering toward equilibrium. This layer-dependent directionality is a direct consequence of the sign rule $\text{sgn}(g_{m,v}^0) = -\text{sgn}(\Delta_{K\Gamma})$ and constitutes a characteristic experimental fingerprint [Fig. 4(a)]. At the optimal gate voltage near the χ_v peak, the transient amplitude can reach several percent of $I_D(\infty)$, growing linearly with ΔV_{GS} for small steps.

Hysteresis under finite-rate gate sweeps.— For a linearly varying gate voltage at rate $r = dV_{GS}/dt$, the re-

laxation equation (3) gives, to first order in $r\tau_{iv}$,

$$f_\Gamma(t) \approx f_\Gamma^{\text{eq}}[V_{GS}(t)] - r\tau_{iv} \chi_v[V_{GS}(t)]. \quad (12)$$

The valley fraction lags the equilibrium value by $\delta f_\Gamma = -r\tau_{iv}\chi_v$. Substituting into Eq. (5), the first-order current deviation is

$$\delta I_D(V_{GS}) = \frac{W}{L} qp V_{DS} \Delta\mu \tau_{iv} r \chi_v(V_{GS}). \quad (13)$$

This approximation is valid for $\tilde{r} \equiv r\tau_{iv} \ll 1$ on the gate-voltage scale over which χ_v varies significantly. The hysteresis current—the difference between forward ($r > 0$) and reverse ($r < 0$) sweeps at the same V_{GS} —is

$$\Delta I_D^{\text{hys}}(V_{GS}) = 2 \frac{W}{L} qp V_{DS} \Delta\mu \tau_{iv} |r| \chi_v(V_{GS}). \quad (14)$$

Equation (14) encodes several signatures that distinguish valley-origin hysteresis from trap-induced effects:

(i) **Profile.** $\Delta I_D^{\text{hys}} \propto qp \chi_v(V_{GS})$ activates only above threshold (where $\chi_v \neq 0$) and tracks the gate-voltage shape of the valley susceptibility. Trap hysteresis follows the trap energy distribution and need not correlate with $\chi_v(V_{GS})$.

(ii) **Layer-dependent transconductance sign reversal.** ΔI_D^{hys} is positive for the bilayer ($\chi_v > 0$) and negative for the trilayer ($\chi_v < 0$). Such a layer-dependent transconductance sign reversal would be difficult to reproduce by conventional extrinsic mechanisms without structural differences between the two devices.

(iii) **Linear sweep-rate dependence.** Valley hysteresis grows as $|r|$, enabling extraction of

$$\tau_{iv} = \frac{|\Delta I_D^{\text{hys}}|}{2|r|(W/L)qp V_{DS} \Delta\mu |\chi_v|} \quad (15)$$

from the slope of peak $|\Delta I_D^{\text{hys}}|$ vs $|r|$, using χ_v obtained independently from a static g_m - V_{GS} measurement [15] or from a first-principles estimate [11, 13]. Trap hysteresis

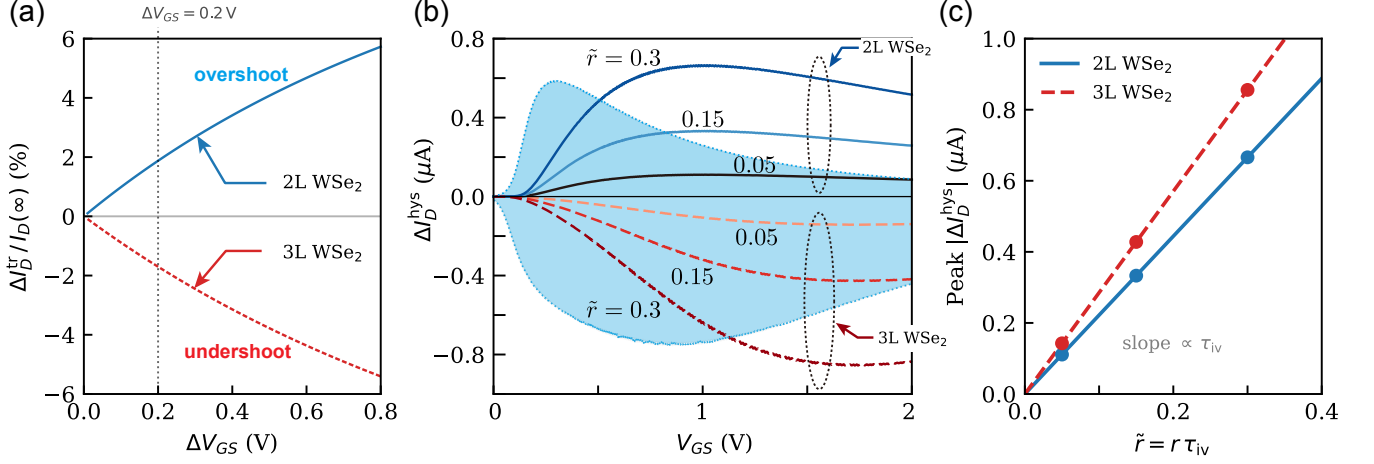


FIG. 4. Low-frequency fingerprints of delayed intervalley relaxation. (a) Signed transient amplitude $\Delta I_D^{\text{tr}}/I_D(\infty)$ vs gate-step size ΔV_{GS} at $T = 300$ K, evaluated at the gate voltage of maximum $|\chi_v|$. Bilayer (blue) overshoots and trilayer (red) undershoots, encoding $\text{sgn}(g_{m,v}^0)$; amplitudes increase linearly before saturating. The dotted line marks the $\Delta V_{GS} = 0.2$ V step of Fig. 2. (b) Hysteresis current $\Delta I_D^{\text{hys}}(V_{GS})$ [Eq. (14)] for bilayer (solid, blue family) and trilayer (dashed, red family) at normalized sweep rates $\tilde{r} = r\tau_{iv} = 0.05, 0.15, 0.30$ (labeled). Opposite signs reflect $\text{sgn}(\chi_v) = \text{sgn}(\Delta K_T)$; grey dotted envelopes trace the normalized shape of $\chi_v(V_{GS})$, confirming the profile correspondence. (c) Peak $|\Delta I_D^{\text{hys}}|$ vs $\tilde{r}$. Linear slopes confirm Eq. (14) and give τ_{iv} directly via Eq. (15).

typically follows a logarithmic or power-law dependence on sweep rate [37–39].

(iv) Subthreshold swing invariance. $\Delta I_D^{\text{hys}} = 0$ in the subthreshold regime where $\chi_v = 0$ [15]. Trap-induced hysteresis is often largest just below threshold where trap filling is most active.

Experimentally, the finite-rate sweep protocol corresponds to a periodic triangular gate-voltage waveform, as commonly employed in cyclic gate-sweep measurements [39]. After initial transients of order τ_{iv} following the onset of the waveform, $f_T(t)$ settles into a reproducible steady-state cycle synchronized to the triangular waveform. In this steady-state limit, each upward and downward branch is locally described by Eq. (12), and Eq. (14) directly characterizes the steady-state hysteresis loop width. This steady-state interpretation renders τ_{iv} extractable from the slope of peak $|\Delta I_D^{\text{hys}}|$ versus $|r|$ without any dependence on initial conditions.

Figures 4(b,c) show the predicted steady-state hysteresis profiles $\Delta I_D^{\text{hys}}(V_{GS})$ for bilayer and trilayer WSe₂ at three normalized sweep rates $\tilde{r} = r\tau_{iv}$, together with the linear scaling of the peak amplitude. The grey envelope traces the normalized shape of $\chi_v(V_{GS})$, confirming the predicted correspondence.

V. TEMPERATURE DEPENDENCE AND SPECTROSCOPY WINDOW

All three signatures are amplified at reduced temperature, and the spectroscopy window shifts into the range of standard laboratory instruments. The valley suscep-

tibility χ_v grows toward the intrinsic bound $(4k_B T)^{-1}$ as T decreases [15], while phonon-mediated intervalley scattering slows as

$$\tau_{iv}(T) = \tau_{iv}^0 [e^{\hbar\omega_{ph}/k_B T} - 1] \quad (16)$$

for a zone-edge phonon with $\hbar\omega_{ph} \approx 25$ meV in WSe₂ [20, 40]. This Bose–Einstein form arises because phonon absorption saturates at low T , leaving stimulated emission as the rate-limiting step [22, 23]. With $\tau_{iv}^0 \approx 60$ ps, this model gives $\tau_{iv}(100 \text{ K}) \approx 1$ ns, placing the spectroscopic peak at $\omega_c/2\pi \approx 160$ MHz—within the bandwidth of commercial lock-in amplifiers and vector network analyzers. Below 50 K, $\omega_c/2\pi$ drops below 30 MHz, and remains accessible with standard instrumentation. For comparison, ultrafast optical spectroscopy reports $\tau_{iv} \sim 0.1$ –1 ps at room temperature in WSe₂ [20, 21, 25], placing $\omega_c/2\pi$ in the 0.1–10 THz range, far above standard electrical bandwidths. The present electrical method becomes accessible when phonon-mediated intervalley scattering slows substantially at reduced temperature, shifting ω_c into the MHz–GHz window mapped in Fig. 5(a), where the shaded band marks the range 10 MHz–10 GHz.

An Arrhenius fit of $\ln \tau_{iv}$ vs $1/T$ within this window yields the phonon coupling energy directly:

$$\frac{d \ln \tau_{iv}}{d(1/T)} \xrightarrow{k_B T \ll \hbar\omega_{ph}} \frac{\hbar\omega_{ph}}{k_B}, \quad (17)$$

providing a purely electrical determination of the zone-edge phonon coupling energy [Fig. 5(b)]. The parameters τ_{iv}^0 and $\hbar\omega_{ph}$ can be extracted from a two-parameter fit to the measured $\omega_c(T)$ curve without any optical access.

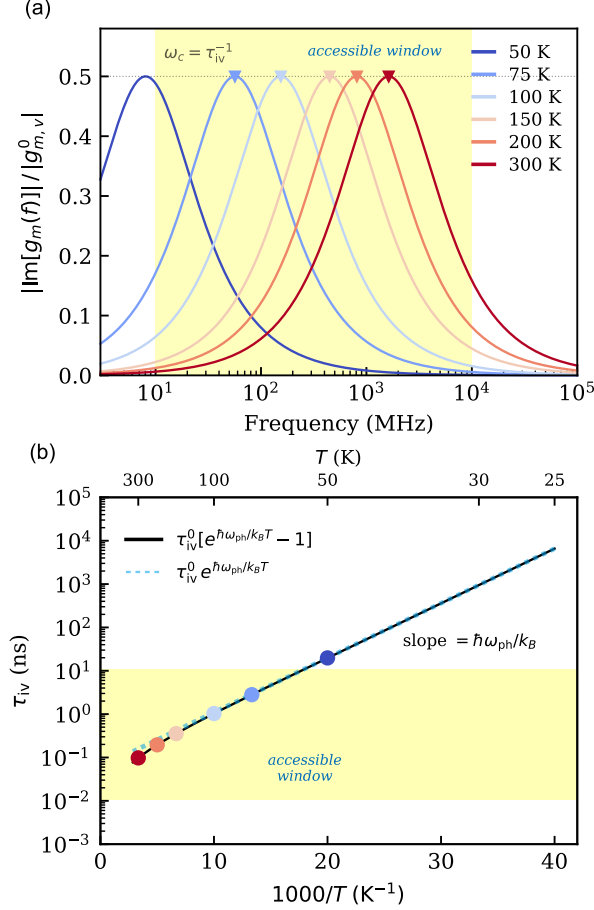


FIG. 5. Transconductance spectroscopy window and Arrhenius analysis. (a) $|\text{Im}[g_m(f)]|/|g_{m,v}^0|$ vs frequency at $T = 50$ – 300 K (blue to red), computed using Eq. (16) with $\hbar\omega_{\text{ph}} = 25$ meV and $\tau_{\text{iv}}^0 = 60$ ps. Triangles mark the characteristic frequency $\omega_c/2\pi$; the shaded band is the accessible window 10 MHz–10 GHz. (b) Arrhenius plot of τ_{iv} vs $1000/T$; the slope of the low-temperature regime yields $\hbar\omega_{\text{ph}}/k_B$ from a purely electrical measurement [Eq. (17)]. Colored dots correspond to the temperatures in (a).

VI. DISCUSSION

These three phenomena provide complementary routes to determine τ_{iv} experimentally. In ac measurements, the frequency sweep at fixed V_{GS}^* gives τ_{iv} from the peak of $|\text{Im}[g_m(\omega)]|$ or the -3 dB rolloff of $\text{Re}[g_m(\omega)] - g_{m,0}$, both pointing to the same ω_c . In transient measurements, the single-exponential decay of $I_D(t) - I_D(\infty)$ directly yields τ_{iv} . In gate-sweep hysteresis measurements, the linear-in- $|r|$ slope via Eq. (15) requires only a static $g_m - V_{GS}$ curve and a rate-dependent cyclic sweep, with no rf electronics.

The quasi-equilibrium limit of Ref. 15 corresponds to $\tau_{\text{iv}} \ll \tau_{\text{tr}}$, where $\tau_{\text{tr}} = L^2/(\mu_K V_{DS})$ is the carrier transit time in the diffusive transport regime. For $W = L = 1 \mu\text{m}$, $V_{DS} = 50$ mV, and $\mu_K = 100 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1}$ [40, 41], the transit frequency is $\tau_{\text{tr}}^{-1} \approx 80$ MHz. If $\tau_{\text{iv}} \sim \tau_{\text{tr}}$, the

crossover in $\text{Re}[g_m(\omega)]$ falls in the 10–100 MHz range, accessible with commercial impedance analyzers. Shorter-channel devices reduce τ_{tr} . In the ballistic limit τ_{tr} is replaced by L/v_{inj} (with v_{inj} the thermal injection velocity), further shifting ω_c to higher frequencies and providing an additional experimental degree of freedom through channel-length scaling of the anomaly. Since τ_{iv} is set by electron–phonon coupling rather than by the channel geometry, it is unchanged by device scaling. Accordingly, all three normalized observables ($\Delta I_D^{\text{tr}}/I_D(\infty)$, $|\text{Im}[g_m]|/|g_{m,v}^0|$, and $\Delta I_D^{\text{hys}}/I_D$) are geometry-invariant. Reducing L at fixed W/L further relaxes the $\tau_{RC} \ll \tau_{\text{iv}}$ condition, because $\tau_{RC} = R_{ch}C_{\text{ox}} \propto L^2$, making sub-micron devices favorable for the measurement. At high frequencies, careful rf de-embedding of contact resistance and pad parasitics will be required to isolate the intrinsic $g_m(\omega)$ from extrinsic contributions. Although the discussion is illustrated for a dual-gate geometry, the mechanism is not specific to a particular gate architecture: gate-all-around devices [14, 42] enhance the electrostatic coupling and hence the magnitude of $p\chi_v$, whereas dual-gate devices offer a more readily accessible experimental platform.

Several other mechanisms can produce frequency-dependent g_m or hysteresis and must be distinguished. Interface states and fast charge exchange are known to degrade the subthreshold swing, whereas border traps produce hysteresis and threshold-voltage drifts [38, 43]. By contrast, the valley hysteresis predicted here arises from delayed intervalley redistribution and does not introduce an additional interface-trap capacitance; within the present model, the subthreshold swing remains unchanged. Charge trapping in dielectrics can produce τ values from microseconds to seconds, but shows no layer-dependent transconductance sign reversal and no dependence on $\Delta_{K\Gamma}$. Applying compressive biaxial strain, which shifts $\Delta_{K\Gamma}$ predictably [14, 44], would modify the valley contribution in a controlled manner, providing a decisive test. Deviations from the predicted Lorentzian line shape of $\text{Im}[g_m(\omega)]$ would directly signal distributed intervalley relaxation channels beyond the minimal single- τ_{iv} model, opening a route to resolving the relaxation spectrum electrically.

VII. CONCLUSION

In summary, we have shown that the valley population in multilayer WSe₂ field-effect transistors acts as a dynamic internal state variable with relaxation time τ_{iv} . The resulting frequency-dependent transconductance $g_m(\omega) = g_{m,0} + g_{m,v}^0/(1 + i\omega\tau_{\text{iv}})$, two-stage transient response, and sweep-rate-proportional hysteresis each provide distinct and electrically accessible probes of intervalley relaxation dynamics. The layer-dependent transconductance sign reversal and gate-voltage profile of all three observables constitute characteristic fingerprints of valley-thermodynamic origin, completing the dynamic

extension of the valley-thermodynamics framework for two-dimensional transistors [15].

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DATA AVAILABILITY

The data used and analyzed during the current study are available from the corresponding author upon reasonable request.

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