Quantum criticality in a uniaxial organic ferroelectric

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Abstract

Tris-sarcosine calcium chloride (TSCC) is a highly uniaxial ferroelectric with a Curie temperature of approximately 130 K. By suppressing ferroelectricity with bromine substitution on the chlorine sites, pure single crystals were tuned through a ferroelectric quantum phase transition. The resulting quantum critical regime was investigated in detail and was found to persist up to temperatures of at least 30–40 K. The nature of long-range dipole interactions in uniaxial materials, which lead to non-analytical terms in the free-energy expansion in the polarization, predict a dielectric susceptibility varying as $1/T^3$ close to the quantum critical point. Rather than this, we find that the dielectric susceptibility varies as $1/T^2$ as expected and observed in better known multi-axial systems. We explain this result by identifying the ultra-weak nature of the dipole moments in the TSCC family of crystals. Interestingly, we observe a shallow minimum in the inverse dielectric function at low temperatures close to the quantum critical point in paraelectric samples that may be attributed to the coupling of quantum polarization and strain fields. Finally, we present results of the heat capacity and electro-caloric effect and explain how the time dependence of the polarization in ferroelectrics and paraelectrics should be considered when making quantitative estimates of temperature changes induced by applied electric fields.

Keywords: quantum criticality, ferroelectric, quantum phase transitions, low temperature, electro-caloric, dipole–dipole interactions, self-consistent phonon theory

(Some figures may appear in colour only in the online journal)

Introduction

There has been a great deal of interest recently in the field of ferroelectric quantum phase transitions [1–3]. A chief reason for this is that the properties of ferroelectrics may be readily tuned with gate voltages and strains, which make quantum ferroelectrics and paraelectrics particularly suited for applications in advanced cryogenic electronics. Highlighting SrTiO$_3$ as an example, one sees that a pristine optically transparent insulator with a band gap of more than 3 eV, and a static dielectric permitivity greater than 10$^5$, may be tuned through an insulator to metal to superconductor transition and back again with the application of just a few volts [4, 5]. On the other hand modest strains [6], chemical doping [7] or isotope substitution [1, 8], can induce ferroelectricity in an otherwise paraelectric ground state. Ferroelectric quantum critical fluctuations have been observed up to temperatures higher than those often seen in other systems, and over a wide range of tuning parameters. Proximity to a displacive ferroelectric quantum critical point where the transverse optical phonon...
frequency becomes very small, and the dielectric function can rise to very high values, is believed to be of importance in understanding superconductivity in materials such as chemically doped [9] or ionic-liquid-gated SrTiO$_3$ [8] and KTaO$_3$ [5], and in oxide interface materials [10] to name just a few.

The nature of quantum criticality in ferroelectrics is strikingly different from that found in other systems, for example magnetic systems [1, 11, 12]. Quantum criticality of the kind considered here arises purely from the atomic vibrations of the lattice and not from electronic or spin degrees of freedom. The excitations around the critical point are propagating bosonic modes (soft transverse-optic phonons), which at the critical point are gapless at the zone centre ($q = 0$). Precessional dynamics and spin flip processes are absent but crucially the free-energy or starting effective action has both short-range and non-analytical long-range interaction terms due to the dipole–dipole interactions present in ferroelectrics. The long range dipole term is Coulombic in origin rather than Amperian (relativistic) as in the magnetic case, and is thus orders of magnitude stronger. In multiaxial systems where the polarization is free to point along the different crystallographic directions, the effects of dipole interactions are relatively benign. Principally they lead to the well known splitting of the frequencies of transverse and longitudinal polar optical phonon modes with respect to the direction of wave propagation $q$ [13, 14]. In a multi-axial displaceable ferroelectric material the frequency of the transverse optical mode disperses as $\Omega^2 = \Delta^2 + v^2 q^2$ with $\Delta$ going to zero at the critical temperature $T = T_C$. The parameter $v$ is the speed of sound of the phonons when the gap $\Delta$ vanishes. The frequencies of the longitudinal optical modes remain finite at the critical point. In the self-consistent field model (alternatively known as the self-consistent phonon model), which has been shown to be quantitatively applicable without any adjustable parameters in a number of ferroelectric systems [1], the correction to the dielectric susceptibility in the paraelectric phase due to quantum critical fluctuations is as explained, as follows more fully in [1, 15–21].

\[
\delta \chi^{-1} \sim \int_0^{\infty} \frac{q^2 n(\Omega_q)}{\Omega_q} dq
\]  

In this equation $\delta \chi^{-1}$ is the change in the inverse electrical susceptibility $\chi^{-1}$ from its zero temperature value, $n(\Omega_q)$ is the Bose function evaluated at the wavevector dependent transverse optical phonon frequency and $q_i$ is a cut-off wavevector typically taken to be the Brillouin zone boundary. For brevity we have neglected the zero point part which is comparatively temperature independent. Note that $\Omega_q = \sqrt{\chi_i^{-1} v^2 q^2}$ where $\chi_i^{-1} = \chi^{-1} + cq^2$ and the constant $c$ is defined below. Therefore equation (1) is a self-consistent equation for $\chi^{-1}$. As explained more fully in [1], it may be solved numerically over the full temperature range from the quantum regime to the classical regime. In the classical regime for a material with or without a finite Curie temperature, and well away from the quantum critical point, the model predicts a Curie–Weiss like susceptibility $\chi^{-1} \sim (T - T_C)$. Close to the quantum critical point the equations become independent of the cut-off wavevector and lead to a temperature dependence which may be expressed in closed form as follows (apart from weak temperature dependent logarithmic corrections associated with marginal dimensionality [1, 16, 18])

\[
\chi^{-1} = a + \frac{5\epsilon_0 k_B^2 T^2}{18hcv}
\]

where $a, b$ and $c$ are the parameters of the Ginzburg–Landau free-energy expansion in the polarization $P$ at zero temperature, i.e. $f = (a/2)P^2 + (b/4)P^4 + (c/2)(\nabla P)^2$, $\epsilon_0$ is the permittivity of free space, $k_B$ the Boltzmann constant and $h$ the Planck constant divided by $2\pi$.

The situation is quite different for the case of a uniaxial system where due to the crystalline details the polarization is confined to vary along one direction only, say the $z$ direction. In this situation as well as the splitting of the frequencies between transverse and longitudinal optical phonon frequencies, the dispersion of the phonons is modified as follows [13, 22]

\[
\Omega^2_q = \Delta^2 + v^2 q^2 + \lambda (\frac{q_z}{q})^2
\]

in which

\[
\lambda^2 = \frac{4\pi Q^2}{\mu V_0}
\]

with $V_0$ the volume of the unit cell, $Q$ the effective charge associated with the soft mode and $\mu$ the reciprocal mass of atoms involved in the vibrations of this mode. The non-analytic character of the last term in the right-hand side of equation (3) (it does not vanish as the wavevector goes to zero) originates from the long-range character of dipole–dipole interactions. This is the case for both classical and quantum critical points. Noting that $(q_z/q)^2 = \cos^2 \theta$ where $\theta$ is the angle between the $z$ axis and the direction of wave propagation $q$, the integral in equation (1) now becomes [16, 23]

\[
\delta \chi^{-1} \sim \int_0^{\infty} \int_0^{\infty} \frac{q^2 n(\Omega_q)}{\Omega_q} dq d\cos \theta
\]  

Again this may be solved numerically giving the full temperature dependence of the susceptibility for uniaxial systems. Effectively the $\cos \theta$ integral acts as one additional dimension and by an appropriate change of variables the integral may be solved analytically close to the quantum critical point where $\Delta$ approaches zero. The result in this case is that the inverse susceptibility varies as the cube of the temperature [16, 23].

\[
\chi^{-1} \sim T^3
\]

Ferroelectricity in TSCC may be tuned to absolute zero using bromine substitution as shown in figure 1. The ferroelectric transition in TSCC is observed to be second order [24] in all experiments carried out so far, an uncommon feature among ferroelectrics, which implies that the zero temperature transition is a quantum critical point (QCP) with quantum
fluctuations persisting over a range of temperatures and doping levels. Independent measurements from several laboratories give the value of the quartic coefficient $B$ in the Landau free energy $G(P, T) = AP^2 + BP^4 + GP^6$ as lying between $+3.1$ to $+4.5 \times 10^{15}$ (SI) [25, 26]; this positive value of $B > 0$ is further evidence that the transition is second-order.

TSCC is a uniaxial ferroelectric with the polarization forming along one direction only (the $b$-axis as indicated in the inset to figure 1). It is highly uniaxial as evidenced for example by its anisotropic electric susceptibility and the values of its lattice parameters—it is pseudo-hexagonal with $(b - c)/c \gg 1$% [30]. It is one of the most prototypical displacive ferroelectrics known [31], in the sense that there exists an under-damped soft mode that can be followed into the low GHz frequency regime [32] from high-$T$ values of $\varepsilon = 630$ GHz $= 21$ cm$^{-1}$. Deuteration [26] produces little change in $T_C$ implying that the transition is not controlled by the N–Cl–hydrogen bonds, compatible with the soft-mode mechanism operating in TSCC [33]. The Curie constant of pure TSCC is less than a hundred Kelvin, the smallest known Curie constant of any known ferroelectric. By way of comparison, typical values for oxide perovskites such as oxygen-18 substituted SrTiO$_3$, BaTiO$_3$ and others are around 50 000 K. This puts TSCC in the family of ultra-weak ferroelectrics (having tiny dipole moments and tiny spontaneous polarizations). An unrelated form of weak ferroelectricity arises in certain multiferroics where small polarizations are induced by magneto-electric coupling.

**Dielectric susceptibility**

The results of dielectric susceptibility measurements taken in a range of brominated TSCC single crystals are presented in figure 2. A description of the crystal growth and experimental procedures can be found in a later section.

The figures show classical Curie–Weiss behaviour in (a) for the 10 per cent bromine doped sample. On the paraelectric side, close to the quantum critical point, $T^*$ behaviour is seen in the 100 per cent sample and over a wider temperature range in the 90 per cent sample. The cross-over temperature $T'$ above which the $T^2$ power law changes to a lower power was taken to be the point where the exponent $\gamma$, from $1/\gamma \sim T'$, deviated from 2 by 10%. This was approximately 30 K in the 90% Br sample and 25 K in the 100% Br sample. Interestingly, in the 90% sample, but absent in the 100% sample, we observe a shallow minimum in the inverse susceptibility around 3 K (figure 2d). Similar minima have been seen in other quantum critical ferroelectrics such as SrTiO$_3$ and KTaO$_3$ [1] and has been explained in those cases quantitatively without adjustable parameters as arising from the electro-strictive coupling of the polarization to strain. This comes from the addition of a term $−g\nabla\phi P^2$ in the free-energy where $\nabla\phi$ is the strain and $g$ the electro-strictive coupling constant [1, 16, 20]. This gives rise to a further correction to the inverse susceptibility varying as $−T^2$ at low temperatures in leading order in $T$ which leads to a minimum [16]. The theory predicts that the minimum should be present in paraelectric samples close to the quantum critical point when $T \ll \Delta$, and so could likely be the explanation for the observations presented here.

Since TSCC is a uniaxial system one might have expected the inverse susceptibility to vary as $T^3$ close to the quantum critical point as explained above. Instead of this we observe dependencies in both $\chi^{-1} \sim T^2$ and $T_C \sim (1 - \lambda x)^{3/2}$ with $\gamma = 2$ as expected and observed in several multi-axial ferroelectrics close to quantum criticality. We can understand this result by noting that the magnitudes of the dipole moments in TSCC are extremely small and thus the effect of the dipole interaction term is greatly attenuated. This is evidenced by a number of observations in TSCC. Firstly, the splitting of the frequencies of transverse (TO) and longitudinal (LO) modes in TSCC is very small (~50 GHz) and as the TO mode approaches zero at the Curie point the LO mode also drops and traces the temperature dependence of the TO mode (but remains finite) [33–35]. The magnitude of the spontaneous polarization in all the ferroelectric samples [36], and the peak value of the dielectric permittivity in all the samples measured, is very small (less than 10). Moreover the peak value gets progressively smaller as the level of bromine doping increases and as $T_C$ is progressively suppressed towards zero [28]. As mentioned in the introduction, the Curie constant ($C$) in TSCC is less than 100 K and is the smallest of all the known ferroelectrics highlighting the fact that the dipole moments are very small in magnitude. Along with the peak value of the electric susceptibility, the Curie constant gets progressively smaller (approximately linearly), as a function of Br doping. This means that the magnitudes of the dipole moments which started off as being ultra-weak in pure TSCC, get progressively
weaker still as one approaches the quantum critical point with Br substitution. This is illustrated in the plot of Curie constant versus Curie temperature in figure 3.

All this essentially means that the $\lambda^2(q^2/|q|^2)$ term in equation (3) is much smaller than $v^2 q^2$ term over the range of wavevectors integrated over in equation (4). This means that the correction to the susceptibility given by equation (4) reduces to the integral equation (1) resulting in an inverse susceptibility varying as $T^2$ as observed, rather than a $T^3$ dependence as expected more generally for uniaxial systems.

Figure 3. The Curie constant $C$ versus Curie temperature in samples of TSCC with varying Br content. The Curie constant was determined from a fit to the part of $1/\chi$ which varied linearly with $T$, i.e. in the classical Curie–Weiss regime of the inverse susceptibility. This held all the way to down to $T_C$ to within one or two degrees for the 0%, 10% and 30% samples and down to the quantum regime, designated by a cross-over temperature $T^*$, in the remaining samples. The uncertainty in $C$ of the data points in the plot is 5%. The error in determining $C$, became much larger for the samples with paraelectric ground states due to the tiny magnitude of the dipoles present in them, a characteristic of the TSCC family of materials. The trend observed in the measured data here was also confirmed from published measurements (e.g. [36]) of the spontaneous polarization $P_S$ for different bromine doping levels using the relation $C(T_{C1})/C(T_{C2}) = P_S^1(T_{C1})/P_S^2(T_{C2})$ from equation (8) where $T_{C1}$ and $T_{C2}$ denote two different Curie temperatures.

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The following analysis outlines a simplified model in which $C$ (and thus $p$ and $\lambda^2$) get smaller as $T_C$ is suppressed. In comparison with ferromagnets, where the Curie constant in the Weiss theory is not an independent parameter but is given by

$$ C = N g^2 \mu_B^2 S(S+1)/(3k_B) $$

(6)

where $g$, $\mu_B$, and $k_B$ are the gyromagnetic ratio, Bohr magneton, and Boltzmann constant, in ferroelectrics $C$ is usually determined empirically by the relation

$$ \epsilon = \epsilon_0 + \frac{C}{T - T_C}. $$

(7)

Some authors [21, 36, 37] have defined a ‘bare’ Curie constant

$$ C_0 = N p^2/k_B. $$

(8)
where \( p \) is the dipole moment and \( N \) the number of dipoles and term the ratio \( C/C_t \) the Rhodes–WohlFarth ratio, in analogy with ferromagnets.

In order to describe Rochelle Salt, Kobeko and Kurchatov [38] assumed a rotating rigid dipole model in which the polarizability

\[
\alpha = p^2/(3k_BT),
\]

where each rotating dipole has dipole moment \( p \). The internal field is given by \( E_i = E(\text{applied}) + \beta P \) where \( \beta \) is the relevant Lorentz factor and \( P \) the polarization (dipole moment per unit volume). This leads to the result \( C = \sigma T_c \) where the constant parameter \( \sigma \) depends on \( \beta \) and the system of units in use. The data for \( C(T_c) \) in TSCC/TSCB are shown in figure 3. Hence in TSCC or its Br-isomorphs, \( C \) decreases towards zero as bromination increases to the QCP. This means that the magnitude of the dipole moments is suppressed as one approaches the QCP and the assumption of a finite \( \alpha^2 \) is not possible, meaning that brominated TSCC may behave as an isotropic material asymptotically as \( T_c \) approaches zero [39]. Note that the results here are not sample dependent. Although the data presented are for samples grown at Princeton, checks were made with our earlier work on specimens from Montana (Schmidt) and Germany (Strauße) [27]. Differences in quoted Curie constant \( C \) arise from the background subtraction schemes employed; subtracting a large background value from small divergences can lead to systematic errors. All data in the present work were processed via the same procedure, so that relative values of \( C(T) \) are reliable, although absolute values less so. Since the Curie Law becomes inexact at low temperature due to quantum effects, this procedure only shows that the general model used is applicable.

\[ \text{Heat capacity and electro-caloric effect} \]

The electrocaloric effect is sometimes predicted to diverge as \( T \) approaches zero (especially close to a quantum critical point where the dielectric function may be large at zero temperature). It is therefore interesting to investigate the electrocaloric effect in materials tuned close to a quantum phase transition to investigate prospects of solid-state refrigeration for applications in cryogenic electronics. The prediction of a diverging electrocaloric effect is due to the fact that from the Maxwell relations, in equilibrium, there is an indirect way of determining the temperature cooling \( \Delta T \) [40–43]:

\[
\Delta T = -T \int \frac{1}{C} \frac{dP(T,t)}{dT} dE
\]

and this expression contains a specific heat term (for which the results of our measurements in TSCC are shown in the inset to figure 4) that diverges faster than its linear pre-factor of \( T \). However that assumes that the dielectric function and corresponding polarization \( P \) are large. In the present case they are exceptionally small, and the numerical values of the electrocaloric effect are extremely small, even at low \( T \) (figure 4).

Also of interest is the time dependence of these relations. The estimate in equation (10) is from the Maxwell relations and generally agrees within a factor of 2 or better with direct measurements of \( \Delta T \) [44, 45]. However, it is not exact. It is based upon equilibrium thermodynamics (dubious in the case of relaxors), and it ignores time dependences. \( P(T) \) actually depends upon time for real ferroelectrics due to relaxation processes [46]. After a voltage is applied \( P(T, t) \) decays with time \( t \), on several different time scales, depending upon temperature \( T \). Typically the long-term value of polarization \( P \) is ca. 50 per cent of the short-term value. In some modern test equipment, the remnant polarization \( P_r \) at long times is denoted with a circumflex.

Thus if \( \Delta T \) is measured on a time scale long compared with the relaxation time \( \tau \), care must be taken to use the proper value of \( P \) in equation (10), which may be only half that of the short-time measured electrical value from the \( I(V) \) hysteresis curve. This will generally lead to an overestimate of \( \Delta T \) inferred from the ‘indirect’ method described by equation (10), and by as much as 50 per cent. No published electrocaloric data include this time dependence of the Maxwell relations. To be more self-consistent, \( \Delta T \) must be measured as a function of time \( t \) in applying equation (10), and the time scale must be commensurate with the relaxation time \( P(T, t) \).

\[ \text{Crystal synthesis and experimental methods} \]

Single crystals (approximately 3 mm × 3 mm × 5 mm) of trisarcosine calcium chloride [(CH₃NHCH₂COOH)₃CaCl₂], as well as bromine doped versions [(CH₃NHCH₂COOH)₃CaCl₂₋ₓBrₓ], were grown by slow evaporation of saturated aqueous solutions. Stoichiometric amounts of sarcosine,
anhydrous CaCl$_2$ and CaBr$_2$ were dissolved into water and thoroughly mixed before being allowed to evaporate over a one week period. In order to determine sample purity and homogeneity, energy dispersive x-ray (EDX) analysis was carried out using a FEI Quanta 200 FEG Environmental Scanning Electron Microscope (SEM) on several doped crystals. The samples were found to be homogeneous with bromine to chlorine ratios matching the expected stoichiometries to within 2%.

Heat capacity measurements were made on a Quantum Design PPMS using the thermal-relaxation method. Measurements of the dielectric permittivity and loss tangent were made with an Andeen Hagerling capacitance bridge and a QuadTech impedance analyzer enabling variable frequency measurements in the range of a few Hz to a few MHz. The sample temperature was varied with a pumped He-3 cryostat with miniature coaxial cables between 0.3 K and 300 K and field ranges) close to the quantum critical point due to the ultra-weak nature of the dipoles (tiny dipole moments), and their further attenuation on approaching the quantum critical point, that the system behaves as an isotropic/multiaxial material especially as $T_C$ approaches zero. The search for $\chi^{-1} \sim T^3$ in a uniaxial quantum critical ferroelectric continues, which has so far remained elusive. We found the estimated electro-caloric temperature change to be small (of the order of a few milli-Kelvin over the accessible temperature and field ranges) close to the quantum critical point due to the fact that TSCC is a system with extremely small dipole moments. For these reasons the entropy change on the application of external electric fields is small.

Conclusions
In conclusion we have made detailed measurements and analysis of quantum criticality in an organic and a uniaxial ferroelectric. At first sight one may have expected the long-range dipole interactions, which are of crucial importance in uniaxial systems, to have drastically changed the power law behaviour of the temperature dependence of the inverse electric susceptibility. This is because for such systems, in general, the effective dimension $d_{\text{eff}} = d + z = 4$ for quantum multiaxial systems, where $d = 3$ is the spatial dimension and $z = 1$ is the dynamical exponent [1]. Instead we found in the special case of the TSCC family of materials that due to the ultra-weak nature of the dipoles (tiny dipole moments), and their further attenuation on approaching the quantum critical point, that the system behaves as an isotropic/multi-axial material especially as $T_C$ approaches zero. The search for $\chi^{-1} \sim T^3$ in a uniaxial quantum critical ferroelectric continues, which has so far remained elusive. We found the estimated electro-caloric temperature change to be small (of the order of a few milli-Kelvin over the accessible temperature and field ranges) close to the quantum critical point due to the fact that TSCC is a system with extremely small dipole moments. For these reasons the entropy change on the application of external electric fields is small.

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References
[1] Rowley S E, Spalek L J, Smith R P, Dean M P M, Itoh M, Scott J F, Lonzarich G G and Saxena S S 2014 Nat. Phys. 10 367
[2] Kim J W et al 2014 Nat. Commun. 5
[3] Horiiuchi S, Kobayashi K, Kumai R, Minami N, Kagawa F and Tokura Y 2015 Nat. Commun. 6
[4] Ueno K, Nakamura S, Shimotani H, Ohtomo A, Kimura N, Nojima T, Aoki H, Iwasa Y and Kawasaki M 2008 Nat. Mater. 7 855
[5] Kawasaki M, Ueno K, Nakamura S, Shimotani H, Yuan H T, Kimura N, Nojima T, Aoki H and Iwasa Y 2011 Nat. Nanotechnol. 6 408
[6] Haeni J H et al 2004 Nature 430 758
[7] Wu H and Shew W 2006 Phys. Rev. B 73 094115
[8] Itoh M, Wang R, Inaguma Y, Yamaguchi T, Shao Y J and Nakamura T 1999 Phys. Rev. Lett. 82 3540
[9] Schooley J F, Hosler W R and Cohen M L 1964 Phys. Rev. Lett. 12
[10] Gariglio S, Reyren N, Caviglia A D and Triscone J M 2009 J. Phys.: Condens. Matter 21
[11] Rowley S et al 2010 Phys. Status Solidi b 247 469
[12] Sachdev S 1999 Quantum Phase Transitions (Cambridge: Cambridge University Press) (ppxiv, 353 p., gbA010388 bnb 2578 Subir Sachdev. ill.; 26 cm. Includes bibliographical references and index)
[13] Strukov B A and Levanyuk A P 1998 Ferroelectric Phenomena in Crystals : Physical Foundations (Berlin: Springer ) (ppx, 308 p., gb9735793 GB98-62489 B A Strukov, A P Levanyuk. 24 cm. Includes bibliographical references (p. 2990) and index)
[14] Lines M E and Glass A M 2001 Principles and Applications of Ferroelectrics and Related Materials (Oxford Classic Texts in the Physical Sciences) (Oxford: Clarendon) (ppxiii, 680 p., gbA001041 1497289 by M E Lines and A M Glass. ill.; 24 cm. Originally published: 1977. Includes bibliographical references and index)
[15] Rochester A B 1971 Sov. Phys.—JETP 33 423
[16] Khmel’nitskii D E and Shneerson V L 1971 Sov. Phys. Solid State 13 687
[17] Khmel’nitskii D E and Shneerson V L 1973 Sov. Phys.—JETP 27 164
[18] Roussev R and Millis A J 2003 Phys. Rev. B 67 014105
[19] Das N and Mishra S G 2009 J. Phys.: Condens. Matter 21 095901
[20] Palova L, Chandra P and Coleman P 2009 Phys. Rev. B 79 075101
[21] Tokunaga M 1988 J. Phys. Soc. Japan 57 4275
[22] Larkin A and Khmel’nitskii D 1969 Sov. Phys.—JETP 25 1123
[23] Khmel’nitskii D E 2014 J. Exp. Theor. Theor. Phys. 118 133
[24] Ashida T, Bando S and Kakudo M 1972 Acta Crystallogr. B 28 1560
[25] Bornarel J and Schmidt V H 1981 J. Phys. C: Solid State Phys. 14 2017
[26] Gergs M. K, Michel D, Schlemmbach H and Windsch W 1986 J. Phys. D: Appl. Phys. 19 2431
[27] Jones S P P, Evans D M, Carpenter M A, Redfern S A T, Scott J F, Straube U and Schmidt V H 2011 Phys. Rev. B 83
[28] Lashley J C, Mumus J H D, Echizen M, Ali M N, Rowley S E and Scott J F 2014 Adv. Mater. 26 3860
[29] Windsch W, Braeter H and Milsch B 1983 Ferroelectrics 47 231
[30] Makita Y 1965 J. Phys. Soc. Japan 20 2073
[31] Feldkamp G E, Scott J F and Windsch W 1981 Ferroelectrics 39 1163
[32] Mackeviciute R, Ivanov M, Banys J, Novak N, Kutnjak Z, Wencka M and Scott J F 2013 J. Phys.: Condens. Matter 25
[33] Kozlov G V, Velkov A A, Scott J F, Feldkamp G E and Petzelt J 1983 Phys. Rev. B 28 255
[34] Prokhorova S D, Smolensky G A, Siny I G, Kuzminov E G and Mikvabia V D 1980 Ferroelectrics 25 629
[35] Feldkamp G E, Douglas K, Lavrencic B B and Scott J F 1980 Bull. Am. Phys. Soc. 25 171
[36] Fujimoto S, Yasuda N, Kashiki H, Takagi K and Fujimoto M 1981 Ferroelectrics 39 1139
[37] Tokunaga M 1987 J. Phys. Soc. Japan 56 1653
[38] Kubeko P and Kurchatov I 1930 Z. Phys. 66 192
[39] Scott J. F 1989 J. Phys. Soc. Japan 58 4487
[40] Correia T and Zhang Q 2014 Electrocaloric Materials (Heidelberg: Springer)
[41] Lu S G et al 2010 Appl. Phys. Lett. 97
[42] Lu S G, Rozic B, Zhang Q M, Kutnjak Z, Pirc R, Lin M R, Li X Y and Gorny L 2010 Appl. Phys. Lett. 97
[43] Mischenko A S, Zhang Q, Scott J F, Whatmore R W and Mathur N D 2006 Science 311 1270
[44] Correia T M, Kar-Narayan S, Young J S, Scott J F, Mathur N D, Whatmore R W and Zhang Q 2011 J. Phys. D: Appl. Phys. 44
[45] Kutnjak Z 2014 Electrocaloric Materials ed T Correia and Q Zhang (Heidelberg: Springer) pp 147–82
[46] Schumacher M, Dietz G W and Waser R 1995 Integ. Ferroelectr. 10 231