Thermodynamics of mono and di-vacancies in barium titanate

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(Dated: 21 December 2013)

The thermodynamic and kinetic properties of mono and di-vacancy defects in cubic (para-electric) barium titanate are studied by means of density-functional theory calculations. It is determined which vacancy types prevail for given thermodynamic boundary conditions. The calculations confirm the established picture that vacancies occur in their nominal charge states almost over the entire band gap. For the dominating range of the band gap the di-vacancy binding energies are constant and negative. The system, therefore, strives to achieve a state in which under metal-rich (oxygen-rich) conditions all metal (oxygen) vacancies are bound in di-vacancy clusters. The migration barriers are calculated for mono-vacancies in different charge states. Since oxygen vacancies are found to readily migrate at typical growth temperatures, di-vacancies can be formed at ease. The key results of the present study with respect to the thermodynamic behavior of mono and di-vacancies influence the initial defect distribution in the ferroelectric phases and therefore the conditions for aging.

PACS numbers: 61.72.Ji 71.15.Mb 71.55.-i 77.84.Dy

I. INTRODUCTION

Barium titanate is a prototypical ferroelectric material. At ambient pressure and temperatures above 393 K it assumes a para-electric, cubic perovskite structure. At lower temperatures a tetragonal distortion of the unit cell is observed which gives rise to ferro-electricity. At even lower temperatures further symmetry reductions lead to ferroelectric orthorhombic and rhombohedral phases. The sequence of phase transitions is the result of subtle structural and energetic differences. Probably the most important technological application of barium titanate is in thin-film capacitors. It also serves as an end member in several lead-free ferroelectric alloys and is used in combination with SrTiO$_3$ to obtain tunable RF devices. Because of its technological importance and its standing as a prototypical ferroelectric, barium titanate has been intensively investigated both experimentally and theoretically.

Intrinsic and extrinsic point defects are of special importance in semi-conductors as they control to a great extent the electronic properties of these materials. In case of barium titanate, this has motivated a considerable number of studies which used conductivity and diffusivity measurements to infer information about the point defect equilibria in this material. It is furthermore assumed that oxygen vacancies play a key role in ageing and fatigue of ferroelectric materials by impeding domain wall motion or by acting as local disturbances of the polarization (see e.g., Refs. 12-19). For this reason, oxygen vacancies and their associates have been intensively studied in experimentally (see e.g., Refs. 21-22) as well as theoretically (see e.g., Refs. 23-32).

Experimental investigation of point defect properties, however, are usually restricted in that they either provide bulk information (e.g., conductivity, diffusivity), which corresponds to a macroscopic average over the sample and is therefore defect unspecific, or very localized information (e.g., electron spin resonance, positron annihilation spectroscopy), which is, however, restricted to certain charge states (unpaired spin states) or types of defects (open volumes, vacancies). In general, the correlation between experimental data and specific defects is indirect and therefore potentially ambiguous. Quantum mechanical calculations on the other hand are capable of providing a very detailed picture of individual defects and simultaneously sample a whole variety of both different configurations and charge states. They are therefore a very powerful tool for verifying model assumptions and for providing fundamental insights into the energetics of defects as well as their kinetic and electronic properties.

The purpose of the present work is twofold. First, we aim to verify the defect models which have been employed to explain the experimental observations (in particular conductivity and diffusivity measurements). Second, we explore the formation of di-vacancies which – as argued above – is an important ingredient for understanding the deterioration of ferroelectric switchability. The calculations are carried out for the cubic phase which is the most stable phase at high temperatures at which most ceramics are processed and at which the initial defect distributions are installed.

II. METHODOLOGY

A. Computational setup

Calculations were carried out within density-functional theory (DFT) using the Vienna ab-initio simula-
tation package (VASP)\textsuperscript{39–42} in combination with the projector-augmented wave (PAW) method to represent the ionic cores and core electrons.\textsuperscript{43,44} In order to find the most suitable representation for the exchange-correlation (XC) functional preliminary calculations were performed within the local-density approximation (LDA)\textsuperscript{45–46} as well as the generalized-gradient approximation (GGA) in the Perdew-Burke-Ernzerhof (PBE) parameterization.\textsuperscript{47} We furthermore considered the effect of treating the Ti-3s and Ti-3p electrons as semi-core states. All calculations were carried out with 20 \(k\)-points for Brillouin zone sampling and a plane wave cutoff energy of 500 eV. For each setup the energy-volume curve of cubic barium titanate was computed and the equilibrium lattice constant, cohesive energy, bulk modulus and its pressure derivative were subsequently obtained from a fit to the Birch-Murnaghan equation of state. Finally, the full band diagrams were calculated at the respective zero pressure lattice constant.

The results of these preparatory calculations are compiled in Table I in comparison with experimental and theoretical data from literature. The influence of the Ti-3p and Ti-3s states on the properties included in this comparison is very small. The most significant difference is the increase of the cohesive energies by about 1\% upon inclusion of the Ti-3s and Ti-3p electrons in the valence. Thus, while for some properties the deep Ti-3s and Ti-3p can play a crucial role, in the present context their effect is expected to be small. For the sake of computational efficiency we did, therefore, not include the Ti-3s and Ti-3p electrons in the valence.

The exchange-correlation functional on the other hand has a more pronounced impact. Within the GGA the bulk modulus is reasonably well reproduced, but the lattice constant is overestimated while the opposite applies for the LDA results. These findings are consistent with the results of previous studies.\textsuperscript{34,48,49} In the present study we have decided to employ the LDA.

In order to determine the properties of the phases in Table I we have employed the same computational settings as for BaTiO\(_3\) but varied the number of \(k\)-points for each material such as to ensure a convergence of the total energy better than 1 meV/unit cell. Energy-volume curves were calculated for each of these structures allowing for full internal relaxation.

B. Defect calculations

1. Formation energies

The formation energy \(\Delta E_D^f\) of a defect in charge state \(q\) is given by\textsuperscript{50–53}

\[
\Delta E_D^f = (E_D - E_H) + q(E_{\text{VBM}} + E_F) - \sum_i \Delta n_i \mu_i
\]  

(1)

where \(E_D\) is the total energy of the defective system and \(E_H\) is the total energy of the perfect reference cell. The second term describes the dependence on the Fermi level, \(E_F\), where \(E_{\text{VBM}}\) is the position of the valence band maximum. The variation of the formation energy with the chemical potentials of the constituents is given by the last term. The difference between the number of atoms of type \(i\) in the reference cell with respect to the defective cell is denoted by \(\Delta n_i\). The chemical potential \(\mu_i\) of constituent \(i\) can be rewritten as \(\mu_i = \mu_i^{\text{bulk}} + \Delta \mu_i\) where \(\mu_i^{\text{bulk}}\) denotes the chemical potential of the standard reference state and is equivalent to the cohesive energy per atom (see Sect. III B).

For the defect calculations we employed supercells with 40, 60, 90 and 135 atoms equivalent to 2 \(\times\) 2 \(\times\) 2 to 3 \(\times\) 3 \(\times\) 3 unit cells. The formation energies given in Table I were obtained by extrapolating the data calculated for different supercell sizes to infinite dilution (see below). We considered both mono-vacancies \((V_O, V_{Ba}, V_{Ti})\) as well as nearest-neighbor di-vacancies \((V_{Ba-V_O}, V_{Ti-V_O})\) taking into account a variety of charge states (see Table III). Due to their size, it is reasonable to assume that host metal ion interstitials have very large formation energies. This assumption is supported by analytical potential calculations which have found Frenkel defects to have much higher energies than Schottky defects.\textsuperscript{54} In order to confirm these results we calculated the formation energies of all interstitials and anti-sites in the neutral charge state. Under metal-rich conditions the formation energies for interstitials and anti-sites are at least 3.8 eV larger than for the lowest neutral vacancy. Under oxygen-rich conditions the difference is smaller but even for the most favorable case (O on Ti anti-site at point D in Fig. 1).
TABLE I. Bulk properties of cubic barium titanate as obtained from experiment and first-principles calculations. US-PP: ultrasoft pseudo-potentials; FP-LAPW: full potential-linearized augmented plane waves; TB-LMTO: tight-binding linear muffin-tin orbitals; ASA: atomic sphere approximation; LDA: local-density approximation; GGA: generalized-gradient approximation; PBE: Perdew-Burke-Ernzerhof parameterization of the GGA; $E_c$: cohesive energy (eV/f.u.); $a_0$: lattice constant (Å); $V_0$: equilibrium volume ($Å^3$/f.u.); $B$, $B'$: bulk modulus (GPa) and its pressure derivative; $E_G^{0\rightarrow T}$: direct band gap at Γ-point (eV); $E_{CB}^{0\rightarrow T}$: indirect band gap measured between points R and Γ; $m_r^*$, $m_h^*$: effective electron (hole) mass at the Γ-point along ⟨100⟩ in units of the electron mass.

| Expt. | DFT (Literature) | DFT (This work) |
|-------|------------------|-----------------|
|       | Ref. 34          | Ref. 38         | Ref. 39       | Ref. 40       | Ref. 39       | Ref. 40       | Ref. 39       | Ref. 40       |
|       | US-PP            | FP-LAPW         | FP-LAPW       | TB-LMTO       | GGA            | LDA            | GGA            | ASA            | LDA            |
| $E_c$ | −31.57 $a$       | −37.92          |               |               | −39.62         | −39.98         | −40.06         | −43.84         | −44.10         | −44.11         |
| $a_0$ | 3.992            | 4.006           | 3.9           | 4.0           | 4.038          | 4.037          | 4.038          | 3.957          | 3.953          | 3.955          |
| $V_0$ | 64.28            | 61.0            | 62.5          | 65.2          | 65.84          | 65.79          | 65.84          | 61.98          | 61.79          | 61.86          |
| $B$   | 173$b$           | 168             | 146           | 185 − 189     | 165            | 161            | 160            | 200            | 194            | 193            |
| $B'$  | 4.5$^c$          |                 |               |               | 4.7            | 4.5            | 4.5            | 4.7            | 4.6            | 4.6            |
| $E_G^{0\rightarrow T}$ | 3.27, 3.38$^d$ | 1.8             | 1.85 − 1.9    | 1.2$^e$       | 1.69           | 1.81           | 1.85           | 1.68           | 1.80           | 1.82           |
| $E_{CB}^{0\rightarrow T}$ |               |                 |               |               | 1.56           | 1.66           | 1.69           | 1.62           | 1.71           | 1.73           |
| $m_r^*$ |                 | 1.16            | 1.16          | 1.16          | 1.10           | 1.00           | 1.01           | 12.8           | 12.7           |                |
| $m_h^*$ |                 | 16.7            | 15.8          | 15.6          | 13.5           | 12.8           | 12.7           |                |                |                |
| $\langle m_r^* \rangle$ |                 | 0.96            | 0.97          | 0.97          | 0.89           | 0.89           | 0.89           |                |                |                |
| $\langle m_h^* \rangle$ |                 | 3.24            | 3.31          | 3.32          | 2.88           | 2.93           | 2.94           |                |                |                |

$^a$ as cited in Reference 32
$^b$ as cited in Reference 33
$^c$ fitted to Rose’s equation of state
$^d$ Reference 35 obtained from optical measurements on tetragonal barium titanate at room temperature.
$^e$ at the experimental lattice constant

the energy difference is at least 0.9 eV. In the following we therefore neglect interstitials and anti-sites and focus exclusively on vacancy defects. It should be pointed out that both metal and oxygen interstitials can play important roles in several other oxides with more open lattice structures (e.g., TiO₂, ZnO, SnO₂). In the case of perovskite lattices the interstitial sites are, however, much smaller and interstitials lead to significant strains which cause large formation energies.

The Brillouin zone integrations were carried out using a non-shifted $2 \times 2 \times 2$ Γ-centered mesh which depending on the symmetry of the defect conﬁguration is equivalent to 4 to 6 $k$-points in the irreducible wedge of the Brillouin zone. The plane wave cutoff–energy was set to 500 eV.

DFT calculations of point defect formation energies in semi–conductors are subject to certain shortcomings, which in order to obtain physically meaningful results must be properly taken into account. Although the band gap is typically underestimated, energy differences within the valence band and conduction bands, respectively, are usually rather well described. A simple correction of the band structure is, therefore, obtained by rigidly shifting the valence band ($\Delta E_{VB}$) and the conduction band ($\Delta E_{CB}$) with respect to each other. The correction energy is then given by

$$\Delta E_{bg} = n_e \Delta E_{CB} + n_h \Delta E_{VB}. \quad (2)$$

For a given defect and charge state the number of electrons in the conduction band $n_e$ was determined by integrating the number of occupied conduction band states. The number of holes in the valence band $n_h$ was similarly obtained as the number of empty valence band states. It should be noted that this correction only considers the effect of the band gap error on the band energy and assumes rigid levels. It does not take into account electronic relaxations which occur if self-interaction effects are properly included.

Since experimentally the band gap of cubic BaTiO₃ cannot be determined for $T \to 0$ K due to the ferroelectric phase transition, we estimated its value by extrapolation of the data at higher temperatures which yields $E_{G}^{expt} = 3.4$ eV (compare Fig. 10 in Ref. 38). In order to be able to correct for the underestimation of the band gap, we furthermore assumed the offset of the calculated band structure to be restricted to the conduction band, i.e., $\Delta E_{VB} = 0$ and $\Delta E_{CB} = E_{G}^{expt} - E_{G}^{calc}.$

Due to the use of supercells elastic and electrostatic interactions between the periodic images of the defects need to be taken into account. Elastic interactions scale inversely with the volume, $O(V^{-1})$, and therefore, have been corrected using finite-size scaling. For charged defects electrostatic interactions due to the net charge moment are present. The corresponding correction term can be expressed in the form of a multipole expansion. The leading term, which describes monopole–monopole interactions, scales with $V^{-1/3}$ and can be calculated explicitly. The next term is due to monopole-quadrupole interactions. It scales with $V^{-2}$ and can in principle also
be evaluated explicitly. Since it displays the same scaling behavior as the elastic interactions, it is, however, conveniently corrected using the same finite-size scaling procedure. In addition, the latter approach has the advantage to avoid ambiguities in the calculation of the moments of the net charge distribution. It also allows to assess the reliability of the correction by means of the error of the linear extrapolation of the data. For strongly delocalized excess charge distributions higher order terms might become important which has, however, not been observed in the present calculations.

Since the magnitude of the electrostatic interactions in condensed matter is screened, the monopole-monopole correction term requires knowledge of the static dielectric constant. For reasons of consistency the latter should be taken from first-principles calculations as well. Using a similar approach as in the present work Ghosez et al. obtained an average value of $\varepsilon = 57$ for barium titanate which was used in the present work.
alent to $2 \times 2 \times 2$ unit cells. Brillouin zone integrations were carried out using the same $k$-point grids as for the formation energy calculations and the plane wave cutoff-energy was again 500 eV. In order to obtain the saddle points, we applied the climbing image nudged elastic band (CI-NEB) method which imposes a minimal number of constraints on the transition path. Only jumps between nearest neighbor sites were considered.

Unlike formation energies migration energies are obtained by calculating differences between configurations which are structurally and electronically very similar. As discussed in detail in Ref. the calculation of migration barriers is, therefore, much less sensitive to the errors described in the following section. For this reason, comparably small supercells should be sufficient and due to the similar electronic structure band gap and potential corrections need not be applied.

III. RESULTS

A. Band structure

In agreement with full potential-linearized augmented plane wave and pseudopotential plane wave calculations our results predict the valence band maximum (VBM) to be located at the R-point. On the other hand, in previous tight-binding linear muffin-tin orbitals calculations within the atomic-sphere approximation the VBM was located at the $\Gamma$-point. In all of these calculations the difference between the highest occupied levels at R and $\Gamma$ was calculated to be on the order of 0.1 eV or less. The effective masses of electrons and holes at the $\Gamma$-point, which can be used for instance for a self-consistent determination of the Fermi level and the charge carrier concentrations are given in Table I.

B. Chemical potentials and stability diagram

In order to be able to derive defect formation energies from total energy calculations, the thermodynamic reservoirs need to be defined. This requires knowledge of the cohesive energies of the constituents in their most stable conformation. Furthermore, as the stability range of barium titanate is restricted by the formation of competing phases such as TiO$_2$ or BaO, the formation energies of these compounds have to be computed as well (see Sect. B1). The results of our calculations of the fully relaxed structures are compiled in Table I in comparison with experimental data. The overall agreement with the reference data is good. In particular, the calculated formation energies compare well with experimental data.

The range within which the chemical potentials of Ba, Ti, and O can vary is restricted by the condition

$$\Delta \mu_{\text{Ba}} + \Delta \mu_{\text{Ti}} + 3\Delta \mu_{\text{O}} = \Delta H_f[\text{BaTiO}_3],$$

which determines the outer triangle in the phase diagram shown in Fig. 1. Considering the compounds given in Table II the chemical potentials are furthermore subject to the following constraints

$$\begin{align*}
\Delta \mu_{\text{Ba}} + \Delta \mu_{\text{O}} &\leq \Delta H_f[\text{BaO}], \\
\Delta \mu_{\text{Ba}} + 2\Delta \mu_{\text{O}} &\leq \Delta H_f[\text{BaO}_2], \\
\Delta \mu_{\text{Ti}} + 2\Delta \mu_{\text{O}} &\leq \Delta H_f[\text{TiO}_2].
\end{align*}$$

They confine the stability range of BaTiO$_3$ to the gray shaded area in Fig. 1. Repeating this analysis with the experimental data gives a phase diagram in good agreement with the calculated one.

C. Defect formation energies

The results of our calculations for the formation energies are summarized in Table II which shows the formation energies of mono and di-vacancies for a Fermi level at the valence band maximum ($E_F = 0$ eV in equation (1)) and for four representative combinations of chemical potentials, which are indicated in Fig. 1. The variation with the Fermi level is shown for the same four cases in Fig. 2. Due to the large formation enthalpy of BaTiO$_3$ the formation energies vary strongly between the extremal points of the phase diagram. In the metal-rich limit (along A–B) oxygen vacancies prevail. They have comparably small formation energies and therefore should be abundant defects. In both cases the formation energies becomes negative for some Fermi level which determines the so-called pinning energy, $\varepsilon_{\text{pin}}$. This implies that under equilibrium conditions the material cannot assume a Fermi level which is closer to the valence band maximum than $\varepsilon_{\text{pin}}$. In the oxygen-rich limit either barium (point C) or titanium vacancies (point D) dominate. Fermi level pinning now occurs in the vicinity of the conduction band, which implies that the Fermi level cannot be pushed arbitrarily close to the conduction band minimum.

The equilibrium defect transition levels can be deduced from the formation energies using equation (4). They are presented in an effective band scheme in Fig. 3. Since vacancies occur in their nominal charge states ($V_{\text{Ti}}^{\text{m}}, V_{\text{Ba}}^{\text{m}}, V_{\text{O}}^{\text{m}}$) almost over the entire band gap, only the band edges are shown. With regard to the effect of the band gap correction term given by equation (2), it turns out that if no band gap corrections are applied the donor transition levels ($q > 0$) are near the calculated conduction band edge, whereas they are near the experimental conduction band edge if the corrections are included. This consistency indicates that the application of the band gap correction described in Ref. is reasonable in the present case.

The location of the levels near the band edges is in agreement with several models which have been developed to reproduce the experimental data (see e.g., Refs. 8–12). It is evident from the extrapolation errors in Table II that typically DFT calculations cannot predict transition levels with an accuracy better than about
FIG. 2. (Color online) Variation of defect formation energies with Fermi level for representative thermodynamic conditions indicated in Fig. 1. The numbers indicate the charge states. Parallel lines correspond to identical charge states. The solid and dashed lines correspond to mono and di-vacancies, respectively. The arrows indicate the position of the Fermi level pinning energy under different conditions.

FIG. 3. Transition levels for mono and di-vacancies in BaTiO$_3$. Only the band edges are shown. The dashed transition levels are positioned inside the valence or the conduction bands (indicated by the light gray shaded areas) and are only included for illustration. The dark grey shaded areas indicate the sum of the extrapolation errors for each transition.

FIG. 4. (Color online) Binding energies for $V_{\text{Ba}}$–$V_\Omega$ and $V_{\text{Ti}}$–$V_\Omega$ di-vacancies as a function of Fermi level. The kinks correspond to charge transition points of the isolated defects (compare Fig. 2 and Fig. 3).

0.1 eV. However, even taking into account this limitation the present results show clearly the intrinsic defect levels to be very close to the band edges ($\lesssim 0.4$ eV). In particular, this applies for the titanium vacancy for which, based on more simplistic calculations, the defect level has been previously calculated to be located $-1.14$ eV below the conduction band maximum. A finding which is confirmed by the present results.

The binding energy of a di-vacancy is given by the difference between its formation energy and the formation energies of the isolated mono-vacancies. It is therefore independent of the chemical potentials. Since charge transitions occur only within about 0.4 eV of the band edges (see Fig. 3), the Fermi level position has very little influence (see Fig. 4). Over the widest range of the band gap the binding energy of the $V_{\text{Ti}}$–$V_\Omega$-di-vacancy is $E_b = -1.93$ eV while the value for the $V_{\text{Ba}}$–$V_\Omega$-di-vacancy is $E_b = -0.62$ eV. Only near the band edges the attraction are somewhat reduced. Thus, the association of metal and oxygen vacancies is energetically strongly favored, especially in the case of the $V_{\text{Ti}}$–$V_\Omega$-di-vacancy.

Note that the formation energy of the $V_{\text{Ba}}$–$V_\Omega$-di-vacancy is constant along lines which are parallel to A–D in the phase diagram (Fig. 1) corresponding to the reaction $\text{BaTiO}_3 + V_{\text{Ba}} - V_\Omega \rightarrow \text{TiO}_2$. Equivalently, along lines parallel to B–D the sum of the formation energies of oxygen vacancies and $V_{\text{Ti}}$–$V_\Omega$-di-vacancies is constant, ($\text{BaTiO}_3 + V_{\text{Ti}} - V_\Omega + V_\Omega \rightarrow \text{BaO}$).
TABLE IV. Calculated migration energies of mono-vacancies in units of eV. The temperature ranges above which the defects become mobile are given in the last column. The negative charge states of the titanium vacancy were not considered since already the neutral charge state displays a huge barrier and, following the trends for the barium and oxygen vacancies, the addition of electrons can only be expected to further increase this value.

| Defect | Charge | Barrier | Onset of mobility |
|--------|--------|---------|------------------|
| $V_{Ba}$ | 0      | 5.82    | > 2500 K         |
|        | −1     | 5.96    | > 2600 K         |
|        | −2     | 6.00    | > 2600 K         |
| $V_{Ti}$ | 0      | 9.84    | > 4300 K         |
| $V_{O}$ | 0      | 1.12    | 490–590 K        |
|        | +1     | 0.97    | 420–510 K        |
|        | +2     | 0.89    | 390–480 K        |

D. Migration energies

The calculated migration barriers are compiled in Table IV. The smallest migration energies are obtained for oxygen vacancies and display a weak charge state dependence. Experimentally the migration barrier for oxygen vacancies has been determined to be $\Delta H_m = 0.91$ eV (Ref. 29). The calculations compare well with this value. In particular, the barrier for the doubly positive charge state, in which the oxygen vacancy should occur for a Fermi level in the middle of the band gap, is in very good agreement with this reference value. We point out that the migration barriers for the different charge states of the oxygen vacancy are very similar to the values for cubic lead titanate calculated by Parke with DFT.28

Using the Einstein relation $6D\tau = \langle r^2 \rangle$, one can estimate the temperature above which a defect becomes mobile by determining the temperature for which $6D\tau$ exceeds $\sqrt{\langle r^2 \rangle_{min}}$ (compare Ref. 62 and 68). The pre-factor for the defect diffusivity can be approximated by the lowest optical phonon mode at the $\Gamma$-point which gives $\Gamma_0 \approx 5$ THz (Ref. 3). If one assumes a typical isochronal annealing time of $\tau = 10$ min and a mean defect separation between $\sqrt{\langle r^2 \rangle_{min}} = 100$ nm and 1000 nm, one arrives at the values which are given in the last column of Table IV. Obviously the only defects, which are fully mobile at typical processing temperatures, are oxygen vacancies. In contrast, due to their very large activation barriers, the migration of metal vacancies is much lower even at temperatures close to the melting point ($\sim 1900$ K, Ref. 65).

IV. DISCUSSION

Experimentally, at low oxygen partial pressures (n-type region) and high temperatures ($\sim 1300 – 1500$ K) the dependence between the electrical conductivity and the oxygen partial pressure is found to be $p_{O_2}^{-1/6}$ which has been assigned to doubly charged oxygen vacancies.8–12 In accordance the present calculations predict oxygen vacancies to be by far the most important defect under metal-rich conditions (low oxygen partial pressure) and to occur in charge state 2+ almost over the entire band gap. At somewhat lower temperatures ($\lesssim 1300$ K) a transition to a $p_{O_2}^{-1/4}$ dependence is observed. At least two different explanations have been discussed in the literature.28 48 Either (1) the charge state of the oxygen vacancies changes from 2+ to 1+ or (2) accidental acceptor dopants are present in the material. In order for the first explanation to be valid, the 2+ /1+ transition level of the oxygen vacancy should be located 1.3 eV below the conduction band minimum.28 However, since the present calculations locate this transition just 0.1 eV below the conduction band minimum, they provide support for accidental acceptor doping as the cause for the change in slope. In fact, a more detailed investigation of the relation between the conductivity and the oxygen partial pressure shows that this mechanism can also explain the $p_{O_2}^{-1/4}$ dependence observed for higher oxygen partial pressures ($p_{O_2} \gtrsim 10^{-2}$ atm, p-type region)48.

The structural and energetic differences between the para-electric, cubic phase and the ferroelectric phases play a crucial role in determining the magnitude and the temperature dependence of ferro-electricity. Since BaTiO$_3$ ceramics are typically processed at temperatures above the cubic–tetragonal phase transition ($T_c = 393$ K) and since the onset of mobility ranges given in Table IV exceed this temperature, the point defect equilibria which are established during cooling should correspond to the cubic phase. In particular, initially (prior to ageing, see Refs. 13 and 10) oxygen vacancies should be rather randomly distributed over the symmetrically inequivalent oxygen lattice sites, even if a clear energetic preference exists for a particular lattice site (which is for instance the case for “c-site” vacancies in tetragonal lead titanate).23,69 In addition, due to the large energy barriers for metal vacancy migration the distribution of barium and titanium, which is established during growth, is expected to be largely maintained if the material undergoes ferroelectric phase transitions.

The binding energies for di-vacancies are negative and large which implies a strong chemical driving force for their formation. Since the migration barriers for metal vacancies are large, they are virtually immobile. In contrast, oxygen vacancies are very mobile at temperatures $\gtrsim 500$ K. Thus, formation of di-vacancies should occur readily at typical growth temperatures by metal vacancies “capturing” diffusing oxygen vacancies.

Defect complexes such as di-vacancies or impurity-vacancy associates carry a dipole moment (see e.g., Ref. 54 for a quantitative calculation). On a cubic lattice different orientations of these defect dipoles are energetically degenerate. In the presence of an electric field15,16 or for non-cubic lattices20 this degeneracy is, however, lifted (i.e., the energies for di-vacancy pairs e.g., oriented along the [001] and [100] axes differs). As argued
above, depending on the barriers, the transition temperature and the cooling rate defect dipoles might not have enough time to achieve the orientation with the lowest energy. Since metal vacancies are rather immobile, re-orientation of these dipoles is much more likely to occur by oxygen vacancy jumps. The gradual reorientation of defect dipoles is then determined by the barrier for oxygen vacancy jumps in the first neighborhood of metal vacancies.

V. CONCLUSIONS

The thermodynamics and kinetics of vacancy defects in paraelectric cubic barium titanate have been studied by means of density functional theory calculations. First, formation, binding and migration energies were derived properly accounting for the shortcomings of the DFT method. The binding energies of metal-oxygen di-vacancies are found to be negative implying that in equilibrium under metal-rich (oxygen-rich) conditions all metal (oxygen) vacancies are bound in di-vacancy clusters. While metal vacancies are practically immobile at realistic conditions, oxygen vacancies can readily migrate at typical growth temperatures. Di-vacancies can, therefore, form if metal vacancies capture one oxygen vacancy.

It is furthermore confirmed that mono-vacancies occur in their nominal (ionic) charge states (\(V_{\text{Ti}}^{m}, V_{\text{Ba}}^{n}, V_{\text{O}}^{p}\)) over the widest range of the band gap. The dominant charge states for the di-vacancies are (\(V_{\text{Ti}}^{m} - V_{\text{O}}^{p}\)) and (\(V_{\text{Ba}}^{m} - V_{\text{O}}^{p}\)). Intrinsic defect levels are confined to a region within \(\sim 0.4\) eV of the band edges. This is partially at variance with earlier calculations based on a more simplistic model, which suggested titanium vacancies to act as hole traps.

The temperatures at which defects are immobilized are higher than or similar to the cubic–tetragonal phase transition temperature. The principal findings of the present study with respect to the thermodynamic behavior of mono and di-vacancies are therefore also relevant for the initial defect distribution in the ferroelectric phases.

ACKNOWLEDGMENTS

This project was funded by the Sonderforschungsbereich 595 “Fatigue in functional materials” of the Deutsche Forschungsgemeinschaft.

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