2 CdZnTe for Gamma- and X-Ray Applications

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2.1 INTRODUCTION

Commercially, cadmium zinc telluride (CdZnTe; CZT) and cadmium telluride (CdTe) are currently the most relevant II–VI compound semiconductor systems for x-ray and gamma-ray radiation detection. Their main advantages over traditional semiconductor materials such as silicon and germanium are their high radiation stopping power due to the larger atomic numbers of their constituents and the low background density of free charge carriers at room temperature (RT) due to their wider bandgap [1,2]. Successful RT spectroscopy and x-ray applications utilizing
CZT or CdTe detectors have been reported and have been in commercial use for many years.

2.1.1 Crystal Growth

Various melt, solvent, and vapor phase crystal growth techniques have been employed over the years to produce CdTe and CZT single crystals [3]. Figure 2.1 shows an 8” CZT ingot grown from a melt using the high-pressure electrodynamic gradient method (HP-EDG) [4]. This growth technique is typically carried out in semiopen growth systems, which can limit the control and reproducibility of the process in a production setting.

A closed-system approach to growing radiation detector-grade CZT is the traveling heater method (THM), in which the CZT crystallizes out of a solution. This is typically done in sealed quartz ampoules at temperatures well below the CZT melting point. Figure 2.2 shows a 3” CZT ingot grown using the THM technique. Figure 2.3 displays an optical photograph of the cross section of the ingot from Figure 2.2 together with an electron backscatter diffraction (EBSD) orientation map of the adjacent slice proving that single crystallinity can be achieved over the entire ingot diameter.

FIGURE 2.1 Typical 8” (24 kg) CZT ingot grown by the HP-EDG technique at eV Products, Inc. (2011).

FIGURE 2.2 Typical 3” CZT ingot grown by the THM technique at eV Products, Inc. (2013). Solvent zone and top (last to crystallize) slice have been taken off from the left. This ingot weighs about 6.3 kg.
Materials and Device Challenges

Significant progress has been made in CZT crystal growth over the past two decades; however, a major hurdle for the broader commercial use of these materials remains the technical challenge of growing larger single crystals, which causes a relatively high cost of large monolithic volume detectors. Furthermore, charge trapping within single crystal grains can cause significant charge loss and, ultimately, poor resolution in gamma spectroscopy. In commercial semi-insulating CZT, the mobility–lifetime products of the holes are typically orders of magnitude lower than those of the electrons \(10^{-6} – 10^{-4}\) cm\(^2\)/V versus \(10^{-4} – 10^{-2}\) cm\(^2\)/V. This difference leads to more or less significant hole tailing in parallel-plate detectors.

This issue is addressed primarily by electron collecting device configurations and correction schemes such as semihemispherical detectors and their variants [5–7], coplanar grid [8], orthogonal and drift strip [9–11], small pixel monolithic array [12,13], and virtual Frisch-grid detector devices [14].

In practice, device design optimization can become rather involved due to the contributions of materials and processing effects originating from nonuniformities of both matrix charge transport properties and electric field distributions, macroscopic defects, electrode interface, and side and street surface properties that can impact the charge collection and noise performance of the final detector devices. In dependence on the actual fabrication process, the interface and surface properties can become more or less dependent on the crystal bulk properties and the crystallographic orientation, which can also affect the electric field distribution within the device. Furthermore, the detector always operates in connection with a particular read-out electronics chain, which opens a matrix of working points that needs to be optimized in dependence on the device design and the materials and interface/surface properties.

While uniform charge loss can be corrected by a variety of techniques, it can still cause difficulties if excessive space charge buildup distorts the internal electric field. In this case, the carrier drift velocities can become position dependent, which can lower the efficiency of, for example, depth-of-interaction-based correction schemes. High-resolution coplanar grid detectors (CPGs) [8,15,16], for example, theoretically require
only uniform (and in practice, also high [17]) electron transport. Good hole transport is in first-order theory not a requirement for high-resolution CPGs. However, if the hole trapping becomes so strong that space charge buildup causes significant local variations of the electric field, the device performance can still be negatively impacted. This space charge buildup occurs upon biasing the detector and deteriorates the device performance already under low-flux conditions (dark polarization). The internal electric field distribution can be affected by the design of the contact barriers [18,19].

This mechanism of space charge buildup in the dark must not be confused with dynamic high-flux polarization, which is triggered primarily by the x-ray or infrared (IR)-generated nonequilibrium carriers and which will be discussed later.

2.1.3 Higher-Dimensional Defects

Nonuniformities of the charge transport properties including defect-related steering effects within single crystals are typically enhanced by the incorporation of second-phase defects such as tellurium (Te) inclusions and extended dislocation networks around them. These difficulties need to be addressed by advances in the crystal growth techniques. Figure 2.4 shows the effect of an inclusion decorated twin boundary on the $x$–$y$ map of charge-collection efficiency as obtained by alpha scanning [20]. Detailed studies regarding the effect of noncorrelated inclusions, dislocations, and subgrain boundaries on charge transport in CZT have been performed using high-spatial-resolution x-ray response mapping [21–23]. The fundamental role of (zero-dimensional) point defects in electrical compensation and uniform and nonuniform charge trapping will be discussed in Sections 2.2 and 2.3.

2.1.4 Annealing

Te inclusions are generated during crystal growth under Te-rich conditions. Within limits, their size distribution can be controlled by crystal growth parameters, such as temperature gradients at the growth interface and so on [24].

---

CdZnTe for Gamma- and X-Ray Applications

As-grown high-ρ

“Clean” low-ρ(n)

Inclusion anneal

“Clean” high-ρ

CT anneal

**FIGURE 2.5** Annealing sequence consisting of an inclusion anneal and a charge transport (CT) anneal.

In principle, Te inclusions can be relatively easily removed by postgrowth annealing in cadmium (Cd)-rich vapor, which is practical to do on a wafer level. The resulting material appears clean in IR transmission; however, it is typically not immediately clear to what extent additional structural damage has been caused in the process as the Te removal or CdTe microgrowth process may create additional defects. Another issue is that the resulting crystals typically turn low-resistivity n-type after the inclusion anneal and a second annealing step usually has to be added to restore the semi-insulating state. This sequence is shown in Figure 2.5. The IR scans of a typical CZT sample before and after annealing are shown in Figure 2.6. After annealing, no inclusions are visible in IR microscopy, but note that this type of microscopy would not reveal any features less than about 1 μm in size nor those that are transparent to IR light.

Figure 2.7 summarizes the electrical properties and charge transport performance of an inclusion-free (6 × 6 × 3) mm³ parallel-plate detector that was prepared from an inclusion and charge transport annealed CZT slice (two-step sequence). The semi-insulating state and radiation detector grade charge transport properties were restored and are comparable with those before annealing. The electron \( \mu \tau \) as obtained from the standard “Hecht fit” [25] of the bias dependence of charge-collection efficiency (CCE) actually improved in this particular case. However, the benefit of such an annealing approach on charge transport uniformity, for example, for fine-pitch x-ray

**FIGURE 2.6** IR scans of the same sample before and after inclusion annealing.
imaging, still needs to be proven. It may always be preferable to minimize second-phase particle incorporation during crystal growth and initial cool down, provided it is technically feasible.

2.1.5 Radiation Sensor Devices

State-of-the-art semiconductor radiation detectors for gamma- and x-ray photons typically employ a device configuration in which the detector is shaped like a rectangular parallelepiped or cylinder. The cathode is typically formed by a continuous metal film deposited on one planar face of the detector. The cathode may extend to the physical edges of the detector or, in certain configurations, partially or all the way over the side-walls [6,7]. Certain device configurations would also employ cathode strips [9]. The anode is typically laid out as a continuous planar electrode, single pixel or small pixel array [12,13], coplanar grid [8], or strip design [9–11]. Additional electrodes such as guard rings, guard bands, steering grids, and control electrodes [26] may be added to the pattern. Multiple metal layers, diffusion barriers, or other coatings on top of the primary electrical contact may be required to ensure mechanical robustness or facilitate
specific interconnect processes. Bare CZT surfaces (sidewalls and streets between pixels or grids) need to be electrically passivated to minimize leakage currents and noise. Some detector concepts require electrically insulated shielding electrodes [27].

Figure 2.8 shows examples of typical detector and electrode configurations. CPGs for high-resolution spectroscopy, linear array detectors for medium-flux x-ray, and large pitch (>1 mm) spectroscopy detectors are shown. The detector thicknesses in Figure 2.8 range from 3 to 25 mm.

The interface processing and electrode technology can significantly affect the electrical properties of the electrode barriers, which constitute for most CZT devices some form of blocking contact, that is, the high-bias resistance will be lower than expected from an ideal ohmic device of a given bulk resistivity [28]. The most common electrode structure for high-resistivity CZT is a back-to-back blocking configuration, but note that the observed barriers will in general be less pronounced than for, for example, CdTe-In diode structures [29].

Figure 2.9 shows the current–voltage (I–V) characteristics of CZT parallel-plate detectors at 23°C. The blocking–blocking (Pt–CZT–Pt) configuration in the ±800 V bias range appears linear in that scale but the zoom shows a well-pronounced, symmetric barrier kicking at about ±1 V. The CZT bulk resistivity can be obtained from guarded electrodes in the linear range below ±100 mV [28]. The I–V characteristics for an asymmetric electrode configuration are shown for comparison. Note that the dark leakage current in nonblocking polarity far exceeds the current expected from an “ohmic device,” that is, additional carriers are injected at higher bias and the detector functions only in reverse polarity.

Figure 2.10 shows a $^{57}$Co spectrum measured with a simple $(13 \times 13 \times 3)$ mm$^3$ parallel-plate detector biased at 700 V. All major features, including the Fe x-rays at 6.4 eV, are well resolved. Figure 2.11 shows a $^{137}$Cs spectrum from the same detector. Note that this is not a favorable device configuration to obtain high resolution and photopeak efficiency at those high energies.

A simple yet highly efficient detector configuration to obtain higher resolution and photopeak efficiency at 662 keV is the coplanar grid design [8]. Figure 2.12 shows the spectra from a set of CZT detectors configured as CPGs. The resolution is well below 2%. The pulser peak is shown to indicate the electronic noise contribution.

Figure 2.13 shows the combined $^{137}$Cs spectrum obtained with a $2 \times 2$ prototype array of $(7 \times 7 \times 20)$ mm$^3$ position-sensitive virtual Frisch-grid detectors with electronically segmented anodes at Brookhaven National Laboratory (BNL) yielding...
FIGURE 2.9  I–V characteristics of CZT parallel-plate detectors at 23°C: (a) Blocking-blocking configuration in ±800 V bias range. (b) Same curve zoomed into the ±20 V range to show the barrier effect. (c) Low bias range precision measurement (±80 mV) to obtain resistivity (guarded electrode). (d) Asymmetric electrode configuration.

FIGURE 2.10  $^{57}$Co spectrum from a (13 × 13 × 3) mm$^3$ parallel-plate detector biased at 700 V.
about 0.8% resolution at 662 keV [30]. The four 20 mm thick CZT base detectors are bar shaped with parallel-plate electrodes as shown at the very right in Figure 2.8.

Figure 2.14 shows the results of an open x-ray beam photon-counting experiment utilizing a 500 μm pitch, 16×16 pixel monolithic array CZT detector. The count-rate response of all 256 pixels to an x-ray flux ramp and the response nonuniformity histogram are shown [31]. Note that from detector to detector, the total counts (sum over all pixels) are usually preserved within 3%. If the count-rate response is stable over time, the pixel count nonuniformities can be easily normalized/calibrated (uniformity correction).
2.2 BASICS OF SEMI-INSULATING CZT

2.2.1 Semi-Insulating State

The temperature and composition dependence of the CZT bandgap can be approximated by Equation 2.1:

\[ \text{FWHM: } \approx 0.8\% \]

\[ \text{Counts} \]

\[ \text{Channel} \]


FIGURE 2.14 (See color insert) Test result for a typical $16 \times 16$ pixel CdZnTe monolithic detector array. (a) The count rate as a function of the x-ray flux (tube current), which is followed by stability and uniformity tests at a constant maximum tube current. (b) The count distribution at maximum current indicating the response uniformity of the device. (From Szeles, C., Soldner, S.A., Vydrin, S., Graves, J., and Bale, D.S., *IEEE-TNS* 54, 1350, 2007.)
where \( x \) is the zinc (Zn) mole fraction in Cd\(_{1-x}\)Zn\(_x\)Te. The other parameters are listed in Table 2.1.

The theoretical maximum resistivity of a semiconductor at a given temperature \( T \) is

\[
\rho_{\text{max}}(T) = \frac{1}{2e_0 n_i(T) \sqrt{\mu_n(T)\mu_p(T)}}. \tag{2.2}
\]

This derives from \( \rho = 1/e_0(\mu_n n + \mu_p p) \) for the case when the mobility–concentration products of the free electrons and holes are equal:

\[
\mu_n n = \mu_p p. \tag{2.3}
\]

The elementary charge is \( e_0 \) and the intrinsic carrier concentration is \( n_i \). Using the parameters from Table 2.1, we obtain \( \rho_{\text{max}} \sim 1 \times 10^{11} \, \Omega \cdot \text{cm} \) for Cd\(_{1-x}\)Zn\(_x\)Te with \( x = 0.1 \) at 23°C. At this point, the conduction type as measured in an ideal thermoelectric effect experiment would be undefined, that is, theoretically neither n- nor p-type. Because of the higher drift mobility of the free electrons in CZT \( \mu_n > \mu_p \), the Fermi level is below the intrinsic level in this case (\( p > n \)). The free-hole concentration \( p = n_i \sqrt{\mu_n / \mu_p} \) is about \( 6 \times 10^6 \, \text{cm}^{-3} \) at RT and the concentration of the free electrons, \( n_i \), is a further order of magnitude lower.

**TABLE 2.1**

<table>
<thead>
<tr>
<th>Cd(_{1-x})Zn(_x)Te Parameters Used for the Calculations</th>
</tr>
</thead>
<tbody>
<tr>
<td>RT electron drift mobility</td>
</tr>
<tr>
<td>Cd(<em>{0.9})Zn(</em>{0.1})Te (cm(^2)/V·s)</td>
</tr>
<tr>
<td>1000 [32]</td>
</tr>
<tr>
<td>RT hole drift mobility</td>
</tr>
<tr>
<td>Cd(<em>{0.9})Zn(</em>{0.1})Te (cm(^2)/V·s)</td>
</tr>
<tr>
<td>50 [32]</td>
</tr>
<tr>
<td>Electron effective mass</td>
</tr>
<tr>
<td>CdTe (0.91×10(^{-30})kg)</td>
</tr>
<tr>
<td>0.11 [33]</td>
</tr>
<tr>
<td>Hole effective mass CdTe</td>
</tr>
<tr>
<td>(0.91×10(^{-30})kg)</td>
</tr>
<tr>
<td>0.73 [34]</td>
</tr>
<tr>
<td>( E_0 ) (eV)</td>
</tr>
<tr>
<td>1.606 [35]</td>
</tr>
<tr>
<td>( a_1 ) (eV)</td>
</tr>
<tr>
<td>0.38 [28]</td>
</tr>
<tr>
<td>( a_2 ) (eV)</td>
</tr>
<tr>
<td>0.463 [35]</td>
</tr>
<tr>
<td>( a_3 ) (eV/K)</td>
</tr>
<tr>
<td>4.5×10(^{-4}) [28]</td>
</tr>
<tr>
<td>( a_4 ) (K)</td>
</tr>
<tr>
<td>264 [36]</td>
</tr>
</tbody>
</table>

This must not be confused with the following two commonly used transition points:

\[ n = p, \]  

which is the point of complete compensation from a Fermi statistics standpoint. The Fermi level is now at the intrinsic level and the semiconductor is n-type as per Equation 2.3 at both sides of this transition point.

\[ \mu_n^2 n = \mu_p^2 p, \]  

is the point where, for example, the Hall coefficient would change its sign. The Fermi level is even further below the intrinsic level than with Equation 2.3 and the material is p-type on both sides of this transition point. While those distinctions are of minor relevance in low-resistivity semiconductors where the Fermi level is energetically far off any of those transition points and majority carrier approximations can be used, they become important in semi-insulators when compensation models are compared with experimental data.

To actually achieve a semi-insulating state at RT, that is, free carrier concentrations of say \(< 10^7 \text{ cm}^{-3}\), requires effective compensation of electrically active foreign impurities, native point defects, point defect complexes, energy levels in the forbidden gap related to crystallographic discontinuities and lattice distortions, and any combinations of these. In practice, shallow donor and acceptor concentrations cannot be matched that closely by controlled (shallow) counterdoping and defects that produce deep energy levels are required to pin the Fermi level close to the middle of the bandgap. Somewhat loosely distinguishing between shallow levels that would be fully ionized when the Fermi level is close to the middle of the bandgap and deep levels that would be only partially ionized to compensate for the net difference of the concentrations of the shallow levels, the condition to obtain semi-insulating material can be generalized [37]:

\[
\sum_i N_{A_{i}}^{\text{shallow}} \leq \sum_j N_{D_{j}}^{\text{shallow}} + \sum_k N_{D_{k}}^{\text{deep}}
\]

and

\[
\sum_j N_{D_{j}}^{\text{shallow}} \leq \sum_i N_{A_{i}}^{\text{shallow}} + \sum_j N_{A_{j}}^{\text{deep}}.
\]

This simply means that the total concentrations of all shallow acceptors, \(N_{A_{i}}^{\text{shallow}}\), must be smaller than the total concentrations of all shallow and deep donors, \(N_{D_{j}}^{\text{shallow}}\) and \(N_{D_{k}}^{\text{deep}}\), together, and at the same time, the total concentrations of the shallow donors, \(N_{D_{j}}^{\text{shallow}}\), has to be smaller than the total concentrations of all the acceptors, \(N_{A_{i}}^{\text{shallow}}\) and \(N_{A_{j}}^{\text{deep}}\). Condition 2.6 remains fully valid if defects with multiple charge states are individually indexed for each of their energy levels. It also provides a sufficient framework for both deep donor and deep acceptor compensation models. Note, however, that a simple comparison of the total numbers of donors and acceptors does not necessarily predict the conduction type of semi-insulating material as this will also depend on the actual deep-level ionization energies.
The calculation of the ionized defect and free carrier concentrations is based on the solution of the charge neutrality condition:

\[ n + k_m \sum_{m} N_{A_m}^{(-k_m)} = p + k_n \sum_{n} N_{D_n}^{(+k_n)}, \tag{2.7} \]

which states that the equilibrium concentration of all negative charges (ionized acceptors \( N_{A_n} \) and free electrons \( n \)) has to equal the number of all positive charges (ionized donors \( N_{D_n} \) and free holes \( p \)); \( k_m \) and \( k_n \) indicate the actual charge states (1, 2, \ldots). The individual concentrations of each species are given by the Fermi–Dirac statistics [38,39]. As an example, the ionized concentration, \( N_{D_j}^{-} \), of a single donor \( j \) with total concentration, \( N_{D_j} \), and degeneracy, \( g_{D_j} \), is given by the energy difference between its 0/+ transition level, \( E_{D_j} \), and the Fermi energy, \( E_F \):

\[ \frac{N_{D_j}^{-}(T,x)}{N_{D_j}} = \frac{1}{1 + g_{D_j} \exp \left( - \frac{E_{D_j}(T,x) - E_F(T,x)}{k_B T} \right)}, \tag{2.8} \]

where:

- \( k_B \) = Boltzmann constant
- \( T \) = temperature
- \( x \) = Zn mole fraction in Cd\(_{1-x}\)Zn\(_x\)Te

Similar equations apply for ionized acceptors. The degeneracy factors, however, may be different for donors and acceptors depending on the band structure of a given semiconductor [40]. Once the neutrality equation is solved for \( E_F \), all individual ionized fractions including the free carrier concentrations can be obtained and the free electron concentration \( n \) is given by

\[ \frac{n}{N_C}(T,x) = \exp \left( - \frac{E_C(T,x) - E_F(T,x)}{k_B T} \right), \tag{2.9} \]

where \( E_C \) is the conduction band minimum energy and \( N_C \) is the conduction band density of states. Note the general temperature and composition dependencies of the donor and acceptor ionization energies in Equation 2.8, which are of particular importance for the proper treatment of localized (deep) levels in compensation models [37].

Figure 2.15 shows the theoretical temperature dependence of the fundamental Cd\(_{1-x}\)Zn\(_x\)Te band edge energies for a composition of \( x = 0.1 \) and the possible energetic positions of donor levels with respect to a host crystal independent reference level, that is, in absolute energy scale [41]. The valance band edge of pure binary CdTe at 0 K was chosen as the energetic zero point and a linear temperature correction was added:

\[ E_V(T,x) = E_V(0 \text{K}, \text{CdTe}) + \alpha_T T. \tag{2.10} \]
The linear temperature coefficient $\alpha_T$ has been taken from a theoretical work [42]. The dotted (TM) lines depict the strong temperature dependence of the ionization energies of localized, host independent donor levels that could be related to, for example, substitutional transition metal impurities [41]. The level related to TM-i would be an intermediate level in the bandgap at RT and would move into the conduction band continuum at high temperatures. The CO-ND line indicates the possible behavior of a donor state that is primarily buildup from cation orbitals, which could cause the level to somewhat follow the conduction band edge with the temperature. In this case, the donor ionization may be expected to stay roughly constant over the entire temperature range. Similar consideration can be applied to acceptor states [37].

These effects have important implications on electrical compensation models that relate high-temperature defect equilibria to spectroscopic data that are obtained at lower temperatures. They also affect the interpretation of temperature-dependent experiments in that the results will depend on assumptions regarding the actual physical nature of the defects.

Condition 2.6 immediately implies that for wide bandgap semiconductors, a semi-insulating state can theoretically always be obtained by excessive doping with deep donor impurities (if the net shallow level concentration is acceptor) or deep acceptor impurities (if the net shallow level concentration is donor), independent of the order of magnitude of the net shallow doping concentration. While this approach will theoretically always produce more or less “expensive resistors,” it will not necessarily lead to semiconductor radiation detectors because the same deep levels (localized centers) that pin the Fermi level close to the middle of the bandgap would also tend to deteriorate the transport/collection of radiation or particle-generated excess charge carriers by trapping.
2.2.2 Charge Trapping and Detrapping

In steady state, the statistics of trapping, detrapping, and recombination via localized centers is governed by the Shockley–Read–Hall (SRH) theory [43,44]. The basic concept is that each trap has an energy level, \( E_t \), in the bandgap with two possible charge states. This is illustrated in Figure 2.16.

The two possible charge states of the generic trap are labeled \( k \) and \( k-1 \), where \( k \) is the more positive charge state in which the defect is “unoccupied.” In general, the \( k/k-1 \) level can belong to deep donors (+/0), deep acceptors (0/–), or higher-charged transitions such as (2+/+) or (–/2–). The four basic processes are (a) electron capture by the unoccupied defect at rate \( R_{c,n} \); (b) electron emission from the occupied defect at rate \( R_{e,n} \); (c) hole capture by the occupied defect at rate \( R_{c,p} \), which is equivalent to electron transfer from the defect to the valence band; and (d) hole emission from the unoccupied defect at rate \( R_{e,p} \), with its equivalent to electron transfer from the valence band to the defect. In this framework, recombination can be understood as a sequence of electron capture (a) followed by hole capture (c), which is essentially an electron transfer from the conduction to the valence band, which eliminates one conduction band electron and one hole in the valence band.

The SRH theory derives a relation between the capture and emission rates based on the principles of detailed balance in thermal equilibrium, which requires equal rates of capture and emission separately for both electrons and holes \( (R_{c,n} = R_{e,n} \) and \( R_{c,p} = R_{e,p}) \) with

\[
R_{c,n} = \nu_n \sigma_n n N_t (1 - P_t),
\]

\[
R_{e,n} = e_n N_t P_t,
\]

\[
R_{c,p} = \nu_p \sigma_p p N_t P_t,
\]

\[
R_{e,p} = e_p N_t (1 - P_t),
\]

\[
\text{(2.11)}
\]

![FIGURE 2.16 Schematic diagram of the four basic capture and emission paths as explained in the text.](image-url)
where:
- \( n \) and \( p \) = total free electron and hole concentrations
- \( N_i \) = total trap density
- \( P_i \) = probability of the trap being occupied, which, for a nondegenerated ground state, is
  \[
P_i = \frac{1}{1 + \exp\left(\frac{E_i - E_F}{kT}\right)}.
  \]  

(2.12)

This leads to expressions for the electron and hole emission rates, \( e_{n/p} \), in terms of thermal velocities, \( v_{n/p}^{th} \), capture cross sections, \( \sigma_{n/p} \), and band densities of states \( N_{c/v} \) including the well-known SRH attempt-to-escape frequency factors \( \omega_{n/p} = v_{n/p}^{th} \sigma_{n/p} N_{c/v} \):

\[
e_n = \omega_n \exp\left(-\frac{E_c - E_i}{k_BT}\right)
\]

(2.13)

\[
e_p = \omega_p \exp\left(-\frac{E_p - E_i}{k_BT}\right).
\]

(2.13)

These relations are then used to determine the net capture, emission, and recombination rates for steady-state (stationary nonequilibrium) conditions under which the net rates of electron and hole capture have to be equal to ensure constant electron and hole concentrations and constant average defect occupations [43]:

\[
U = R_{e,n} - R_{e,p} = R_{e,p} - R_{e,n}.
\]

(2.14)

With Equations 2.11 and 2.13, the average occupation probabilities, net capture, and net recombination rates can be calculated from Condition 2.14. Those quantities obviously depend on the injection level (\( \Delta n \) and \( \Delta p \)). The net recombination rate \( U \) is

\[
U = \frac{N_i v_{n}^{th} \sigma_n v_{p}^{th} \sigma_p \left(np - n_i^2\right)}{v_{n}^{th} \sigma_n \left[n + n_i \exp\left(\frac{E_i - E_l}{kT}\right)\right] + v_{p}^{th} \sigma_p \left[p + n_i \exp\left(\frac{E_i - E_l}{kT}\right)\right]},
\]

(2.15)

where \( n = n_{eq} + \Delta n \) and \( p = p_{eq} + \Delta p \) denote the total carrier concentrations in the bands and \( E_i \) is the intrinsic Fermi level. Equation 2.15 necessarily implies \( U = 0 \) for thermal equilibrium (\( np = n_i^2 \)).

Note that while Condition 2.14 is always fulfilled in steady state, it is not straightforward to apply the SRH treatment to experiments where significant
occupation transients occur, such as the analysis of emission rates in photoinduced current transient spectroscopy (PICTS) or thermoelectric effect spectroscopy (TEES) [45].

### 2.2.3 Capture Cross Sections

Figure 2.17 shows a schematic illustrating the efficiencies of deep donors and deep acceptors with respect to electron and hole trapping in a fully compensated semiconductor. In the figure, the (+/0) donor and (0/–) acceptor levels are at the same energy in the bandgap. The shallow levels are fully ionized. The terms *weak* versus *strong* are based on simple considerations regarding the Coulomb potential of the trapping center with respect to each respective carrier type [46].

There is no well-established theory regarding the exact mechanism of the non-radiative transition of free carriers from the bands into localized levels close to the middle of the bandgap [47]. Single longitudinal optical phonon energies are about 21.3 meV in CdTe [48], which is obviously too low to facilitate the required energy transfer with the crystal lattice. One commonly considered mechanism is multiple phonon emission mitigated by significant lattice distortions [49]. This is usually conceived by plotting the electronic and elastic energies of the band extrema and defects over the lattice displacement (configuration coordinate diagrams) [50].

Without addressing the possible implications of the energy dissipation problem, the effects of modified phonon fields around defects, or experimental issues with deep-level spectroscopy that can lead to surprisingly low capture cross sections apparent for those defects [47,51,52], the following classic estimates on physical trapping cross sections may be applied to obtain estimates on upper limits for the allowed deep-level defect concentrations in radiation detector grade semi-insulating CZT. To this end, we consider only the trapping efficiencies of deep donors and deep acceptor per Figure 2.17, which compensate the net shallow level doping $N_{AS}^{tot} - N_{DS}^{tot}$.

---

**FIGURE 2.17** Schematic diagram illustrating the efficiencies of deep donor and acceptor levels with respect to electron and hole trapping in a fully compensated semiconductor.
according to Condition 2.6. To calculate the actual ionized fractions of the deep donor and acceptor states, the neutrality equation, Equation 2.7, has to be solved for \( E_F \) as described in Section 2.2.1.

The simple consideration that the attractive potential of neutral point defects should not extend much over the range of the actual lattice distortion [46] puts the effective capture radius in the order of approximately one lattice constant:

\[
    r_{\text{neutral}} \sim \frac{\alpha_{\text{CdTe}}}{2}.
\]  

(2.16)

This gives a capture cross section for the neutral defect of about \( \sigma_{\text{neutral}} \sim 3 \times 10^{-15} \text{ cm}^2 \) for both electron trapping on neutral acceptors and hole trapping on neutral donors. The lattice constant of CdTe is \( a_{\text{CdTe}} \), which is about 6.48 Å at RT [53].

For charged point defects, the effective capture radius may be estimated by comparing the Coulomb energy of the singly charged defect with the thermal energy of the charge carrier [46], which gives

\[
    r_{\text{singly}}(T) \sim \frac{e_0^2}{6\pi\varepsilon_0\varepsilon_r k_B T},
\]

(2.17)

and a capture cross section of \( \sigma_{\text{singly}} \sim 4 \times 10^{-13} \text{ cm}^2 \) at RT for both electron trapping at positively charged donors and hole trapping at negatively charged acceptors.

In this picture, the capture cross sections for electrons and holes at the same defect would differ by two orders of magnitude with \( \sigma_e > \sigma_h \) for deep donors and \( \sigma_h > \sigma_e \) for deep acceptors. Looking at the simplifications of this approach, including physical insufficiencies such as the use of the bulk dielectric constant, \( \varepsilon_r \), to calculate the attraction between point charges, the implied temperature dependence in Equation 2.17 should not be taken too seriously at this point and experimental tests and quantum-mechanical modeling sometimes seem to confirm [49,54] but often show otherwise [55].

### 2.2.4 Deep-Level Concentrations

Using the rough estimates from Equations 2.16 and 2.17, the maximum allowed deep-level concentrations to maintain sufficient charge transport for radiation detector applications can be obtained by a minimum trapping time consideration. The trapping time (often referred to as “lifetime” in this context) is the average time that a free charge carrier can exist in the bands before trapping occurs. The electron lifetime components relate to the capture rates from Equation 2.11, such as

\[
    \tau_e' = \frac{n}{R_{e/n}'} = \frac{1}{\sqrt{\pi} \sigma_{\text{e}} N_{i}'}
\]

(2.18)

for electron trapping at level \( j \) that has an unoccupied concentration of \( N_i' \) (the positively charged donors or the neutral acceptors in Figure 2.17). The capture cross section of the unoccupied trap \( j \) is \( \sigma_{\text{e}} \) for electrons and the thermal velocity of the
electrons is $\nu^e_{th}$ as the directed drift motion induced by an external bias is, in most practical applications, small compared with the random thermal motion. A similar relation can be derived for hole trapping with the only difference being that the concentrations of the occupied defects must be used. Note that the electron or hole effective masses (Table 2.1) should be used to calculate the $\nu^h_{th}$ as the carriers move in the respective bands before they get trapped. For the simple scenario illustrated in Figure 2.17, where each defect in a specific charge state accepts only one carrier type, the total trapping time for, for example, the electrons $\tau_{e}^{tot}$ can be written as

$$\frac{1}{\tau_{e}^{tot}} = \sum_{j=1}^{N} \frac{1}{\tau_{e}^{j}} = \frac{1}{\tau_{e}^{ID}} + \frac{1}{\tau_{e}^{NA}},$$

(2.19)

where $\tau_{e}^{ID}$ and $\tau_{e}^{NA}$ are the ionized deep donor and neutral deep acceptor–related trapping time components, respectively. Using the capture cross sections according to Equations 2.16 and 2.17 and the RT mobility data from Table 2.1, the upper limits for deep-level concentrations in radiation detector grade CZT are obtained, as shown in Table 2.2.

These data are reasonably consistent with experimental data obtained on CdTe:Cl/In and CdTe:Ge/Sn [56,57] and trapped carrier concentrations of $(10^{12}–10^{13})$ cm$^{-3}$ were determined by space charge measurements [19,58].

Note that due to the lack of a developed theory regarding the actual capture/reemission mechanisms from localized deep centers as outlined before, fundamental misjudgments cannot be ruled out. Otherwise, the implications would be significant: If impurity concentrations beyond a few parts per billion (ppb) are detected in the radiation sensor grade CZT using macroscopic mass spectroscopy, these impurities do not generate deep levels (isolated or as complex), are almost completely compensated, or they have an active concentration in the crystal matrix that is much lower than the concentration reported by the mass spectroscopy (e.g., due to second-phase gettering or direct precipitation due to low solubility).

The relevant deep-level defects control the semiconductor’s electric and charge transport properties already at concentrations well below the detection limits of standard mass spectroscopic techniques. These concentrations are also very low

<table>
<thead>
<tr>
<th>TABLE 2.2</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Maximum Deep-Level Concentrations According to Equations 2.18 and 2.19 to Obtain Mobility–Lifetime Products as Indicated</strong></td>
</tr>
<tr>
<td>---</td>
</tr>
<tr>
<td>$\mu \tau &gt; 10^{-3}$ cm$^2$/V</td>
</tr>
<tr>
<td>Ionized deep traps</td>
</tr>
<tr>
<td>Neutral deep traps</td>
</tr>
<tr>
<td>$\mu \tau &gt; 10^{-5}$ cm$^2$/V</td>
</tr>
<tr>
<td>Ionized deep traps</td>
</tr>
<tr>
<td>Neutral deep traps</td>
</tr>
</tbody>
</table>

*Note:* The capture cross sections are from Equations 2.16 and 2.17.
compared with any technically manageable purity of the crystals or intentionally introduced doping concentration levels.

2.2.5 SELF-COMPENSATION

Another fundamental implication comes from the requirement that the ionized deep levels have to compensate the overall net shallow level concentration offsets to achieve a semi-insulating state according to Condition 2.6. The limits per Table 2.2 simply mean that any shallow defect states need to be compensated by other shallow levels of opposite sign within about 0.1 ppb. This is an extremely close compensation requirement given the fact that intentional doping with shallow donor impurities in concentrations of several parts per million (ppm) can be technically necessary to overcompensate the residual impurity background of similar concentration.

Such close compensation can only be achieved by an energetically favorable, self-driven mechanism that compensates or deactivates a wide range of shallow level concentrations without introducing deep defects. Several specific mechanisms and their combinations have been proposed and investigated since the 1950s, mostly in context with the opposite technical problem of achieving high n- or p-type conductivity in the more ionic II–VI and III–V semiconductor systems for optoelectronic applications such as light-emitting devices [59,60].

In the latter case, self-compensation needs to be prevented or at least limited while in the radiation detector field, and it is exactly this mechanism that helps to achieve semi-insulating properties at relatively low deep-level concentrations. While successful qualitative and quantitative microscopic proof, good matching of experiment and model predictions, and full preparative control were achieved in some cases such as for ZnSe [61–63], any specifics regarding the compensation mechanism leading to semi-insulating CZT will not be detailed here given the current state of insufficient experimental evidence.

The common denominator for all possible self-driven electrical compensation mechanisms is the spontaneous formation of compensating defects, such as additional native point defects. This process is Fermi level–driven because the formation enthalpy of the charged defects with respect to the neutral ones is lowered by the energetic difference between the Fermi level and the defect level [40,64]. Under conditions close to the thermodynamic equilibrium and for the example of intentional doping with shallow donor impurities and compensation by (relatively) shallow acceptor-like native defects, the process would drive the Fermi level down to a critical value below which further generation of native defects would no longer be energetically favorable. This “optimized” Fermi level position also depends on the temperature and partial pressures in that the component partial pressures at a given temperature define the chemical potential for the generation of the neutral defects as, for example, vacancies. If the formation enthalpy and the energy levels associated with those native defects are such that the intrinsic Fermi level can be closely approached within the p-T existence region of the compound, semi-insulating material can be obtained, theoretically even without the introduction of deep levels. The spontaneous defect generation stops when the formation enthalpy of the charged native defects approaches zero (negative energies are required to drive the
process). An additional practical requirement is that the corresponding defect equilibriump can be either frozen-in (quenching) or maintained by an appropriate cooling regime down to temperatures where kinetic barriers eventually stop any further uncontrolled equilibration.

The formation enthalpies and final defect states can also be affected by lattice deformations/relaxations or actual configuration changes; for example, calculations suggest that the Te antisite in the Jahn–Teller distorted configuration [65] may generate a donor level close to the middle of the bandgap in CZT [66,67]. Certain impurities can also be incorporated in multiple configurations, for example, substitutional and interstitial. If they act as donor in one configuration and as acceptor in the other configuration, they will be built into either configuration at a Fermi level–driven probability and partially compensate each other (autocompensation) [68,69]. Moreover, complex formation (e.g., A-centers) during cool-down from equilibration temperature can further alter the charge transport properties at RT.

To treat this field with appropriate respect, the reader is referred to the original works of deNobel [70] and Kroeger [71] and some excellent review articles [59,60].

### 2.3 Dynamic High-Flux Polarization

#### 2.3.1 X-ray-Generated Charge

A primary mechanism of the energy dissipation of x-ray photons in semiconductors is multiple electron–hole pair generation by photoelectric excitation. The x-ray energy affects both the number of generated free carriers and the absorption profile.

In the case of full absorption of an incident x-ray flux, $\Phi$, by photoelectric effect only, the pair-generation rate per detector cross section is

$$g_{\text{pair}} \sim \Phi \ast \frac{E_{\text{mean}}}{E_{\text{pair}}} \cdot$$  \hspace{1cm} (2.20)

The mean energy of the incident x-ray energy spectrum is $E_{\text{mean}}$ or, in the case of thinner detectors, the mean energy of the absorbed photons. The average electron–hole pair-creation energy, $E_{\text{pair}}$, in semiconductors is typically about three times the bandgap energy [72], which amounts to about 4.6 eV for Cd$_{1-x}$Zn$_x$Te with $x \sim 0.1$ at RT.

At higher incident x-ray energies, the probability of Compton scattering increases, in which case part of the energy is dissipated by the recoiling electron by subsequent photoelectric excitation of electron–hole pairs while the scattered photon will either dissipate its energy by photoexcitation, undergo another Compton scatter event, or escape from the detector, in which case charge is lost with respect to full energy deposition.

A general description of all possible primary and secondary processes occurring during the interaction of ionizing radiation with semiconductors is rather complex and excellent reviews can be found in the literature [73].

#### 2.3.2 Space Charge and Detrapping

Trapped charge carriers can cause electric field distortions within the detector. Applying the Poisson equation:
to, for example, a constant positive space charge, \( q(r) = q_0 \), shows that a trapped hole density of only \( 1.3 \times 10^{11} \, \text{cm}^{-3} \) would be sufficient to completely shield the electric field at the anode of a 2 mm thick detector biased at 1000 V. Note that the dielectric relaxation time \( \tau_{\text{rel}} \):

\[
\tau_{\text{rel}} = \frac{\varepsilon_0 \varepsilon_r \rho}{\varepsilon_0},
\]

is on the order of (20–30) ms in CZT at RT, which means that a semiconductor can stay charged with respect to the outside circuitry for a significant amount of time.

According to Section 2.2.2, there are basically two ways for a trapped excess charge carrier to be removed from a semiconductor: it can either be re-emitted into the respective band and (possibly after a sequence of multiple trapping/detrapping events) eventually drift out of the device or it can recombine at the localized defect site if the defect captures a charge carrier of the opposite sign. The probability of either event occurring depends on the defect properties (concentration, ionization energy, capture cross section), the temperature, and the available concentrations of nonequilibrium carriers of both types (injection level, i.e., photon or particle flux and energies). At low photon fluxes, the net recombination rates at the deep levels are very low and the dominant process at low photon fluxes is detrapping.

### 2.3.3 Polarization in Photon Counting

The phenomenon, which can be called “ballistic deficit assisted high-flux polarization” under intense x-ray irradiation is caused by the space charge accumulation of photogenerated holes close to the cathode, which leads to electric field pinching [75,76]. Electron charge clouds deposited on the cathode side must travel through the low-field region and therefore slow down. As the flux increases, the transit time can increase dramatically. Due to the ballistic deficit in the counting electronics, x-ray deposition events are recorded at lower energies. The entire spectrum shifts and more and more counts vanish below the counting threshold. This can lead to a sudden decrease in the photon count rate with increasing x-ray flux if a critical flux is exceeded. Independent experimental confirmation regarding the field deformation can be found in, for example, the work by Prekas et al. [77]. Note that in direct current (DC) mode, the steady-state photocurrent will always monotonically increase with flux, even in the high-flux polarization regime. Figure 2.18 shows the concept of electric field pinching due to space charge buildup [75].

### 2.3.4 Dynamic Lateral Lensing

Another effect related to the internal charging of defects is dynamic lateral polarization, which can cause flux-dependent charge steering in high-flux applications [78]. It is caused by the same mechanism that leads to photon-counting polarization and is observed as lateral charge lensing [79].
Figure 2.19 shows the pixel counts from a 500 μm pitch monolithic CZT detector array that polarizes [78]. The open x-ray beam was collimated to a 4 mm diameter area. The thin red circle marks the irradiated area. All pixels respond in a typical way to low flux. With increasing flux, a nonpolarizing device would show increasing intensity in the same area. In the polarizing device, however, counts are steered inward. The outer irradiated pixels progressively lose counts as the flux increases and eventually the whole device shuts off and only a few counts are seen in a donut-shaped pattern [78].

### 2.3.5 X-ray Photocurrents

Without recombination losses, the photocurrent density, \( j_{\text{photo}} \), should approach a bias-independent value on the order of

\[
j_{\text{photo}} \sim e_0 g_{\text{pair}}. \tag{2.23}
\]

![Diagram of electric field pinching](image1)

**FIGURE 2.18** Illustration of electric field pinching due to space charge buildup. (From Bale, D.S. and Szeles, C., *Phys. Rev.*, B 77, 35205, 2008.)

**FIGURE 2.19** (a) Dynamic lateral lensing in a polarizing pixel array detector. (b) Pictorial with collimated flux on a polarizing pixelated detector showing the lensing electric field lines and the lateral component of the electron velocity. Top: the collected charge profile is shown for this situation. (From (a,b) Soldner, S.A., Bale, D.S., and Szeles, C., *Trans. Nucl. Sci.*, 54, 1723, 2007; (b) Bale, D.S., Soldner, S.A., and Szeles, C., *Appl. Phys. Lett.*, 92, 082101, 2008.)
If currents beyond the value per Equation 2.23 are measured, additional carrier injection via the electrodes must have occurred (photo-assisted breakdown).

Figure 2.20 shows dark and steady-state x-ray photo I–V curves from parallel-plate detectors of different CZT ingots. The two flux levels are high enough that the measured total currents at any given bias are much higher than the DCs and approximate the x-ray-generated steady-state photocurrents. The nominal incident x-ray flux densities coincide reasonably well with the numbers estimated from the photocurrents per Equation 2.23.

To understand the minor differences between samples from different ingots, the actual photocurrent transients must be looked at, as more or less severe recombination losses can lead to a decay of the photocurrent $j_{\text{photo}}(t)$ from its initial ($t=0$) value and Figure 2.20 resembles only the steady state.

The polarization dynamics will manifest themselves in the photocurrent transients. After a step-like flux change, it will take a finite time for the defect occupations and recombination-relevant transition rates to settle, that is, the photocurrent stabilizes at the rate with which the carrier system attains dynamic equilibrium with the photon field after flux changes [80]. In the case of sudden exposure, this is seen as a decay of the initial photocurrent toward a lower steady-state value (DC polarization) [80,81]. A similar time delay will be seen if the flux is reduced as the dynamics are simply reversed (“after-glow” if the detector was polarized in the high-flux state).

Figure 2.21 shows typical RT photocurrent transients of CZT detectors with different defect structures. While the electron mobility–lifetime products are similar for both detectors ($\sim 5 \times 10^{-3} \, \text{cm}^2/\text{V}$), the hole transport is significantly better for detector (a).

The photocurrent transients are controlled by the same dynamics that lead to electronically assisted count-rate polarization but there is an important difference
between the two modes: The measured DC photocurrent is independent of the temporal trapping of the photogenerated charge as long as detrapping is more likely than actual charge loss due to recombination. There is no ballistic deficit in current integrating measurements so that a simple slowdown of the transient charge carriers cannot reduce the photocurrent. However, trapping can still indirectly affect the photocurrent as it deteriorates the electric field by space charge buildup (in the same way as described earlier for photon counting), at which point local charge accumulation together with high-generation rates can increase the recombination losses.

At higher bias, recombination is the only mechanism that can reduce the DC photocurrent from the expected levels per Equation 2.23 and DC polarization is therefore expected to occur at much higher fluxes than the (electronically assisted) count-rate polarization. The steady-state photocurrent will always increase monotonically with the flux even if the detector is significantly polarized as increased recombination balances the electric field distortion by removing space charge. The degree of polarization at given operating conditions (flux, bias, temperature, etc.) depends critically on the defect structure of a specific CZT sample.

Independent of the particular defect structure, CZT x-ray detectors can be “prepolarized” by subbandgap IR illumination in which case the photocurrent response to sudden x-ray exposure no longer shows any delay [80]. Figure 2.22 shows the temporal photocurrent response of a low $\mu\tau_h$ detector to a sudden high-flux x-ray exposure: curve (a) has been measured under standard conditions in the dark and curve (b) was obtained when the same detector was “conditioned” by constant IR illumination at a wavelength of 880 nm (~1.41 eV), which excites the entire detector volume. Note, however, that at this point, the internal electric field is massively distorted and the transit time for individual electron clouds traveling through the whole thickness of the device is dramatically increased.

### 2.3.6 IR Photocurrents

For materials characterization experiments aimed at studying DC polarization or preselecting material for high-flux applications, IR illumination can be alternatively...
used instead of x-ray radiation to probe the high-flux response of the detector. This allows for more simple, less expensive, and more efficient setups, especially in terms of fast-flux switching.

Figure 2.23 compares the x-ray and IR-induced photocurrent responses of a CZT detector. Three different x-ray energies have been used, which affects the absorption profile. The steady-state values of all x-ray and IR-induced photocurrents were adjusted to the same value by regulating the x-ray flux and the IR diode current, respectively. The initial decay is more pronounced in the case of the IR measurement.
CdZnTe for Gamma- and X-Ray Applications

for two reasons: In the case of IR illumination, there is no convolution of the fast initial decay with the relatively slow x-ray shutter motion. In fact, the photocurrent follows the nominal diode rise upon switching instantaneously (90% in 1.5 μs). The other difference is in the absorption characteristics. While the mean absorption depth for x-rays increases with the energy (largest decay for 130 kVp in Figure 2.23), IR illumination with energies below the bandgap energy generates electron–hole pairs almost uniformly throughout the entire detector volume.

Another fundamental difference between x-ray and IR excitation is the generation path of the electron–hole pairs. At subbandgap photoexcitation, the electron–hole pair generation can only occur via defect states in the forbidden gap. This has the important implication that the bulk generation efficiency in a certain exposure regime depends on the specifics of the defect structure of the semiconductor, that is, already the initial photocurrent ($t=0$ value before polarization onset) contains information about the defect structure. The photocurrent evolution thus represents a convolution of charge generation and flux-dependent charge loss due to recombination (polarization dynamics) and both the initial photocurrent and the steady-state current depend on the defect structure. This is not the case for x-ray irradiation where the generation rates are simply defined by Equation 2.20. However, in the limit of full charge collection and no higher-order effects, similar generation rates, $g_{pair}$, will produce similar photocurrents.

Figure 2.24 shows the IR flux dependence of the photocurrent transients at constant bias. A notable manifestation of the convoluted dynamics is the appearance of temporal quasi-stable photocurrent plateaus upon flux onset. The time period during which the photocurrent stays stable at its initial, nonpolarized value increases as the flux decreases (delayed polarization).

![Figure 2.24](See color insert) Example of the IR flux dependence of the photocurrent transients at constant bias.
Solid-State Radiation Detectors

Figure 2.25 shows the bias dependence of the photocurrent transients at constant IR flux. At higher bias, the peak photocurrents are essentially bias independent if plateaus develop, that is, if polarization does not kick in instantaneously upon flux onset. A quasi-stable plateau response at higher bias is reached after ~80 μs and the plateaus extend in time as the bias increases (delayed polarization).

The flux and bias dependencies of the nonpolarized initial peak (plateau) and the polarized steady-state currents are summarized in Figure 2.26. As expected, the peak photocurrents are essentially bias independent (if the bias is not too low). At any given bias, both the peak and the steady-state photocurrents increase nearly linearly with the flux so that the ratio between the peak and the steady-state currents becomes essentially flux independent. This implies that at any flux, about the

**FIGURE 2.25** *(See color insert)* Example of the bias dependence of the photocurrent transients at constant IR flux.

**FIGURE 2.26** Peak and steady-state photocurrents from Figure 2.24 compared with the corresponding numbers at one-tenth of the original bias voltage.
same relative fraction of the photogenerated charge is lost due to recombination. This can be understood considering that in these experiments, even the lowest IR flux still generates a charge equivalent of several $10^7$ x-ray photons/(mm$^2$s) at 120 kVp. Eventually, the steady-state photocurrents are naturally strongly bias dependent in polarizing material as the degree of polarization is reduced as the bias increases.

2.4 SUMMARY

This chapter provided a brief introduction to a variety of materials and device aspects of CZT radiation sensors for gamma- and x-ray applications. The basics of semi-insulating semiconductor detector crystals were explained in terms of electrical compensation and charge trapping statistics. The fundamental implications of practically observed charge transport properties in detector grade CZT on the allowed deep-level defect concentrations were discussed in detail and the concept of electrical self-compensation was outlined. The final section focused on high-flux polarization phenomena in both photon counting and photocurrent modes.

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